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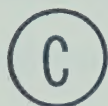
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A STUDY OF THE $^{37}\text{Cl}(p,n)^{37}\text{Ar}$ REACTION

by



ANTHONY CYRIL WANKIEWICZ

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THE UNIVERSITY OF ALBERTA

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The undersigned certify that they have read, and recommend to the Faculty of Graduate Studies for acceptance, a thesis entitled A STUDY OF THE $^{37}\text{Cl}(p,n)^{37}\text{Ar}$ REACTION, submitted by ANTHONY CYRIL WANKIEWICZ, in partial fulfilment of the requirements for the degree of Master of Science.

ABSTRACT

Neutrons from the eight lowest states of Argon-37 have been observed with the $^{37}\text{Cl}(p,n)^{37}\text{Ar}$ reaction using the time-of-flight technique. Using a 17 keV thick NaCl target, enriched in ^{37}Cl , neutron angular distributions were obtained for the ground, 1.41 and 1.61 MeV levels of ^{37}Ar , at mean proton energies of 4.95, 5.00 and 5.05 MeV. The angular distributions show a rapid variation of shape with energy. Further neutron angular distributions, using a target of ethyl chloride, 150 keV thick, were obtained at 5.22 MeV for the ground, 1.41, 1.61, 2.22 and 2.489 MeV levels of ^{37}Ar . The distributions using the thicker target were also somewhat asymmetric about 90° suggesting a possible direct interaction component. An attempt was made to compare these cross sections with those predicted by the statistical theory. The spins and parities of all lower levels except the 2.22 MeV ^{37}Ar state are known. Since the comparison of cross sections normalized to the ground state is very poor for those levels whose spins are known it was not possible to assign a spin to the 2.22 MeV level by this method. The absolute cross section of the ground state is only 20% of the statistical prediction.

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INTRODUCTION

The purpose of the experiment was to investigate the nucleus of Argon-37 through the $^{37}\text{Cl}(p,n)^{37}\text{Ar}$ reaction. Extensive studies of its levels up to 2.8 MeV have been made but the spin and parity of the 2.22 MeV level has not been uniquely determined. This was attempted in the present work by comparing the experimental cross section of the $^{37}\text{Cl}(p,n)$ reaction to the 2.22 MeV level in Argon-37, with the prediction of Hauser-Feshbach theory.

This work is divided into four chapters. The first chapter summarizes the Hauser-Feshbach or statistical theory for describing compound nuclear reactions. There the differential cross section is interpreted as decay into competing channels following the formation of the compound nucleus. The branching ratios of this decay are determined by the transmission functions for penetration of the optical model potential barriers for each channel. Chapter two summarizes previous work on the lower levels of Argon-38. It then goes on to describe the compound nucleus Argon-38, first to consider whether its level density at the excitation energy of this experiment is high enough to warrant a statistical interpretation to the cross sections and then to review work on the excitation function for its formation.

Chapter three describes the accelerator and time-of-flight system used to obtain neutron spectra from both

the solid and gaseous targets. The last chapter presents the experimental neutron angular distributions along with an extensive critique of random and systematic error. The results of the theoretical calculation are compared with experimental cross sections and a conclusion given. Some details of the theory and gas cell construction are found in the appendices.

CHAPTER ONE

THEORY OF COMPOUND NUCLEAR CROSS SECTIONS

1.1 Cross Sections from the Collision Matrix

Nuclear interactions are usually classified as either direct interactions or compound nuclear interactions. The $^{37}\text{Cl}(p,n)^{37}\text{Ar}$ reaction does not contain a strong direct component at the energy used in this study. The following discussion deals mainly with those reactions proceeding through a compound nucleus though much of it applies to direct or prompt reactions as well. The point of departure will be exclusive use of the Bohr assumption, or alternatively, the exclusive use of compound nucleus levels of energy very near to that of the system energy.

The background to compound nucleus reaction cross sections contains many approaches⁽¹⁾, all of which involve approximations to the effects of a large number of compound nucleus resonances and channels. We can regard these channels as different forms of the wave function ψ depending on which pair of particles are regarded as independent. If ϕ is the original plane wave aimed by the experiment at the target, then $\psi = S\phi$, where S is the evolution operator for this transition⁽²⁾. Since this is a steady state interaction, energy will be conserved so that S will

have only the elements ⁽²⁾ $U_{nm} \delta(E_n - E_m)$ where E_n, E_m are the kinetic energies of the channels n and m . S and U are referred to as the scattering and collision matrices respectively.

U has many useful properties. First, for a rotationally invariant interaction the U_{nm} are symmetric, i.e. $U_{nm} = U_{mn}$ ⁽²⁾. Secondly, it gives the transition probability, P_{nm} , between plane waves in different channels through ⁽²⁾:

$$(1) \quad P_{nm} = \frac{1}{2\pi\hbar} \delta(E_n - E_m) |U_{nm}|^2 \quad \text{for } m \rightarrow n ,$$

where $\hbar$ is Planck's constant over 2π . To get the differential cross section for this transition we must first integrate $P_{nm} \rho(E_n)$ over the continuum of final states n , where $\rho(E_n)$ is the density of final states in a unit volume per solid angle. Then this total transition rate is divided by the incident flux.

If $U_{nm} = \delta(n-m)$ defines no collision at all, the cross section must be proportional to $|U_{nm} - \delta(n-m)|^2$ since we can only measure the elastically scattered wave, not the total wave in the entrance channel. The result is usually written in a form which transforms the collision matrix elements to explicit values of ℓ , i.e. of each incident -- orbital angular momentum -- partial wave.

The angular dependence is then contained in the spherical harmonics. The differential cross section for spinless particles is then ⁽²⁾:

$$(2) \quad \frac{d\sigma_{ba}}{d\Omega} = \frac{\pi}{k_a^2} \left| \sum_{\ell} (2\ell+1)^{\frac{1}{2}} i^{\ell} Y_{\ell 0}(\theta, \phi) (U_{nm} - \delta(n-m)) \right|^2.$$

k is the channel wave number and $Y_{\ell 0}(\theta, \phi)$ a spherical harmonic over the spherical coordinates θ and ϕ . Here (n, m) are more specific than (b, a) since the former refers to specific values of ℓ (and later to channel spin s and total angular momentum J also) for each pair of emergent and incident particles, b and a respectively, with b leaving the residual nucleus in a given excited state. If we introduce channel spin s , the total angular momentum, J , is no longer just equal to ℓ . We get in fact ⁽¹⁾:

$$(3) \quad \frac{d\sigma_{ba}}{d\Omega} = \frac{\pi}{k_a^2} \sum_{ss'} \left\{ \begin{matrix} \ell & s & m_s \\ \ell' & s' & m_{\ell}' & m_s' \end{matrix} \right\} \frac{1}{(2I+1)(2i+1)} \left| \sum_{J\ell\ell'} (2\ell+1)^{\frac{1}{2}} \right. \\ \times (\ell s 0 m_s | J m_J) (\ell' s' m_{\ell}' m_s' | J m_J) \\ \left. \times i^{\ell'} Y_{\ell' m_{\ell}'}(\theta, \phi) (U_{nm} - \delta(n-m)) \right|^2$$

where once again, the U_{nm} are based on momenta, and the two quantities in bra-ket notation are Clebsch-Gordan coefficients. The z -components of the various angular momenta are labelled m and the incident particle and target

intrinsic spins are I and i respectively. The formula is for unpolarized beam and target as can be seen from the sum over channel spins being outside the norm. m refers to the set (a, ℓ, s, m_s) and n to (b, ℓ', s', m_s') .

The interference between different partial waves disappears when we integrate the cross section over all angles. The result:

$$(4) \quad \sigma_{ba} = \frac{\pi}{k_a^2} \sum_{ss\ell} \frac{(2J+1)}{(2I+1)(2i+1)} |U_{nm} - \delta(n-m)|^2.$$

Since $U_{nm} = U_{mn}$, the reciprocity theorem follows by equating this equation term by term.

$$(5) \quad (2I_a+1)(2i_a+1)k_a^2 \sigma_{ba}^J = (2I_b+1)(2i_b+1)k_b^2 \sigma_{ab}^J.$$

We have shown explicit dependence on J for further application below. It is convenient here to define the total reaction cross section to be $\sum_b \sigma_{ba}$ with the elastic channel excluded. We then use the unitarity property of the collision matrix to write⁽¹⁾:

$$(6) \quad \sigma_{\text{reac}}(a) = \frac{\pi}{k_a^2} \sum_{Js\ell} \frac{(2J+1)}{(2I+1)(2i+1)} \{1 - |U_{mm}|^2\}$$

in terms of diagonal matrix elements only.

1.2 Average Cross Sections Over an Energy Interval

Experimentally, cross sections show rapid fluctuations with energy^(3,4). These have been explained as interference between many narrow but overlapping resonances in the compound nucleus continuum. Consequently we next consider cross sections averaged over an energy interval much larger than these resonances. The resulting equation for the differential cross section is further simplified if we assume that the cross-terms in the expansion of equation 3 can be ignored. This is the random phase approximation i.e. the matrix elements have a rapid but independent energy variation with a mean of zero. The result is usually converted to a form which explicitly shows the angular dependence⁽¹⁾. Thus for $a \neq b$:

$$\begin{aligned}
 (7) \quad \overline{\frac{d\sigma_{ba}}{d\Omega}} &= \frac{\pi}{k_a^2} \sum_{ss's_s m_s} \frac{1}{(2I+1)(2i+1)} \sum_{J\ell\ell'} (2\ell+1) \\
 &\times (\ell s 0 m_s | J m_J)^2 (\ell' s' m_{\ell'} m_{s'} | J m_J)^2 \\
 &\times \overline{|U_{nm}|^2} \sum_L \left(\frac{2\ell'+1}{4\pi} \right) (\ell\ell' 00 | L0) (\ell\ell' 0 -m_{\ell'} | L0) \\
 &\times P_L(\cos \theta) \quad ,
 \end{aligned}$$

where $P_L(\cos \theta)$ are Legendre polynomials, and the horizontal bar signifies the mean value. Sums over these coefficients, called Z-coefficients, are found in standard tables^(5,6).

The average integral cross section is exact, however.

From equation 4 we have simply:

$$(8) \quad \bar{\sigma}_{ba} = \frac{\pi}{k_a^2} \sum_{ss'J\ell\ell'} \frac{(2J+1)}{(2I+1)(2i+1)} \overline{|U_{nm}|^2} \quad a \neq b .$$

It can be shown that for scattering ($a = b$), the energy averaging removes the interference of the potential scattering from the compound elastic scattering integral cross section⁽⁸⁾. If the latter is added to the average of equation 6 we arrive at the following expression for compound nucleus formation:

$$(9) \quad \bar{\sigma}_{\text{comp}}(a) = \frac{\pi}{k_a^2} \sum_{JS\ell} \frac{(2J+1)}{(2I+1)(2i+1)} \{1 - |\bar{U}_{mm}|^2\} .$$

The first factor in each term can be thought of as the 'size' of each partial wave plus the probability of appropriate orientation of the momenta to give J , while the second factor is the probability of penetrating the potential barriers and forming the compound nucleus. This second factor is called the transmission function for that channel.

$$(10) \quad T_m \equiv 1 - |\bar{U}_{mm}|^2 ,$$

where $m = (a, J, s, \ell)$, indicates the channel.

1.3 Bohr Assumption

The Bohr assumption expresses the independence of formation and decay of the compound nucleus.

$$(11) \quad \bar{\sigma}_{ba}^J = \left(\frac{\pi}{k_a^2} \sum_{s\ell} \frac{(2J+1)}{(2I+1)(2i+1)} T_m \right) \left(\frac{G_{bJ}}{\sum_c G_{cJ}} \right) \quad a \neq b ,$$

the second term being the branching ratio. If we now substitute these expressions for $\bar{\sigma}_{ba}^J$ and $\bar{\sigma}_{ab}^J$ into the reciprocity theorem, equation 5, we get

$$(12) \quad \frac{G_a^J}{G_b^J} = \frac{\sum_{s\ell} T_m}{\sum_{s'\ell'} T_n} .$$

Since the G's are not yet normalized we are free to set:

$$(13) \quad G_a^J = \sum_{s\ell} T_a .$$

Thus

$$(14) \quad \bar{\sigma}_{ba} = \sum_J \bar{\sigma}_{ba}^J = \frac{\pi}{k_a^2} \sum_{\substack{Js\ell \\ s'\ell'}} \frac{(2J+1)}{(2I+1)(2i+1)} \frac{T_m T_n}{\sum_{cs''\ell''} T_p}$$

$$a \neq b$$

$$p = (c, J, s'', \ell'')$$

follows from substituting equation 13 into equation 11.

Comparing this equation with equation 8 gives the identity:

$$(15) \quad \overline{|U_{nm}|^2} = \frac{T_m T_n}{\sum_{CS''\ell''} T_p} \quad a \neq b .$$

Substituting this expression into equation 7 relates the differential cross section to the transmission coefficients.

$$(16) \quad \frac{d\sigma_{ba}}{d\Omega} = \frac{\pi}{k_a^2} \sum_{ss'm_s m_s'} \frac{1}{(2I+1)(2i+1)} \sum_{J\ell\ell'} (2\ell+1) \\ \times (\ell s m_s | J m_J)^2 (\ell' s' m_{\ell'} m_{s'} | J m_J)^2 \\ \times \frac{T_m T_n}{\sum_{CS''\ell''} T_p} \sum_L \left(\frac{2\ell'+1}{4\pi} \right) (\ell\ell' 00 | L0) (\ell\ell' 0 -m_{\ell'} | L0) \\ \times P_L(\cos \theta) , \quad a \neq b .$$

If the residual nucleus is left excited in a continuum of states then the $\overline{d\sigma_{ba}/d\Omega}$ will be the cross section for leaving the residual nucleus in the energy range E , $E+dE$ if we multiply equation 16 by $\rho^J(E)dE$, where $\rho^J(E)$ is the level density. Also, the denominator, $\sum_{CS''\ell''} T_p$ must be taken to be a sum over the discrete states and an integral of $T_p \rho^J(E')dE'$ over the continuum states. The $\rho^J(E)$ follows from the total level density $\rho(E)$ plus the spin distribution assumed.

1.4 Transmission Functions

We now require the determination of the T_m , or rather, the $|\bar{U}_{mm}|^2$, from which they are defined. (We could also

talk of an averaged complex phase shift defined by $\bar{U}_{mm} = \exp 2i\delta_m$). These diagonal matrix elements are completely determined by the wave function in the elastic scattering channel. The T_m are usually expressed in terms of the logarithmic derivative of the wave function at R' beyond which the nuclear part of the force disappears. The wave function or its logarithmic derivative are, in turn, subject to the coulomb, centrifugal and nuclear potentials through the Schrödinger equation.

The form of this potential is:

$$(17) \quad V(r) = V_{\text{coulomb}} + V_{\text{centrifugal}} + V_{\text{nuclear}}$$

where r is the separation distance of the particle of reduced mass μ , from a nucleus of charge Ze and atomic number A . V_{coulomb} will be Ze/r outside the nucleus and a function of the charge density distribution within. $V_{\text{centrifugal}}$ is equal to $\ell(\ell+1)\hbar^2/2\mu r^2$.

The optical model has the general form for V_{nuclear} of:

$$(18) \quad V_{\text{nuclear}} = (-V_0 - c_0 E - iW_0) \xi_0(r) + i(W_1 + c_1 E) \xi_1(r) + V_{\text{spin orbit}}$$

The first term is the volume part of the potential and the second term the surface part. V_0 is the real refractive well depth; E , the incident energy and both $+iW_0$ and $+iW_1$ the imaginary or absorptive well depths. c_0 and c_1 are

first derivatives of these quantities with energy. $\xi_0(r)$ and $\xi_1(r)$ are form factors. $\xi_0(r)$ is usually of the Woods-Saxon form:

$$\xi_0(r) = (1 + e^{(r-R_0 A^{1/3})/a_0})^{-1}.$$

$R_0 A^{1/3}$ and a_0 are called the real well radius and diffuseness respectively. $\xi_1(r)$ can be the derivative of a Woods-Saxon form, or a Gaussian form:

$$\xi_1(r) = \exp(-(r-R_1 A^{1/3})^2/a_1),$$

where $R_1 A^{1/3}$ and a_1 are the imaginary well radius and diffuseness. In this work, we take the former shape and W_1 is set equal to $4a_1 W_2$ by convention. $V_{\text{spin-orbit}}$ is the spin-orbit potential:

$$(18A) \quad V_{\text{spin-orbit}} = \left(\frac{\hbar}{m_\pi c}\right)^2 V_{\text{so}} \frac{1}{r} \frac{d}{dr} \xi_0(r) \vec{L} \cdot \vec{\sigma}$$

where m_π is the pion mass and $\vec{L}$ and $\vec{\sigma}$ are the orbital and spin angular momenta operators. 'c' is the speed of light.

Such a well exhibits single particle resonances -- the size or giant resonances which are often mirrored in the transmission functions. The expressions for T_m are involved, so an approximate result for only one size resonance at $E = E_0$ and no spin-orbit potential is shown below⁽⁹⁾:

$$(19) \quad T_m \approx (4P_m \gamma_m^2 W') / \{ (E_O - E - S_m \gamma_m^2)^2 + P_m \gamma_m^2 + W'^2 \} .$$

P_m is the penetrability of the potential barriers.

$$P_m \approx (k_a R' / \{ F_m^2(k_a R') + G_m^2(k_a R') \}) (\pi K a_O \coth \pi K a_O)$$

where F_m and G_m are regular and irregular solutions to the radial Schrödinger equation outside the nuclear force and $K = \frac{1}{\hbar} \sqrt{2\mu V_O}$.

γ_m^2 is the single particle reduced width, here equal to $\hbar^2 / \mu R'^2$ and W' is the average of the absorptive potential over the region where the real part of the potential is attractive. S_m is the shift function and equals

$$- \ell + \left\{ \frac{r F_m dF_m/dr + r G_m dG_m/dr}{F_m^2 + G_m^2} \right\} \Big|_{r=R'} .$$

Size resonances occur at energies

$$E_m \sim \frac{\hbar^2 \pi^2}{2\mu R'^2} \left(\lambda + \frac{1}{2} \right)^2 .$$

(λ an integer) above the bottom of the well and their natural width in a real well is $\sim 2P_m \gamma_m^2$.

The transmission functions are close to 1 for large penetrabilities but when $T_m \ll 1$, they reflect these size resonances. This can be seen in a further approximation to equation 19⁽¹⁰⁾.

$$(20) \quad T_m \approx 4P_m \{ \gamma_m^2 W_O' / ((E_O - E)^2 + W_O'^2) \} \quad \text{for } T_m \ll 1 .$$

The ratio in parentheses is called the strength function. Both equations apply when the effect of distant resonances is small. See Vogt⁽¹⁾ for more general formulae. If all permissible exit channels and their properties are known and their transmission functions determined, equation 16 can be used to calculate differential cross sections to any one channel. It can be shown that the sum of transmission functions over all exit channels can be expressed in terms of compound nucleus level widths and densities (see Appendix 1).

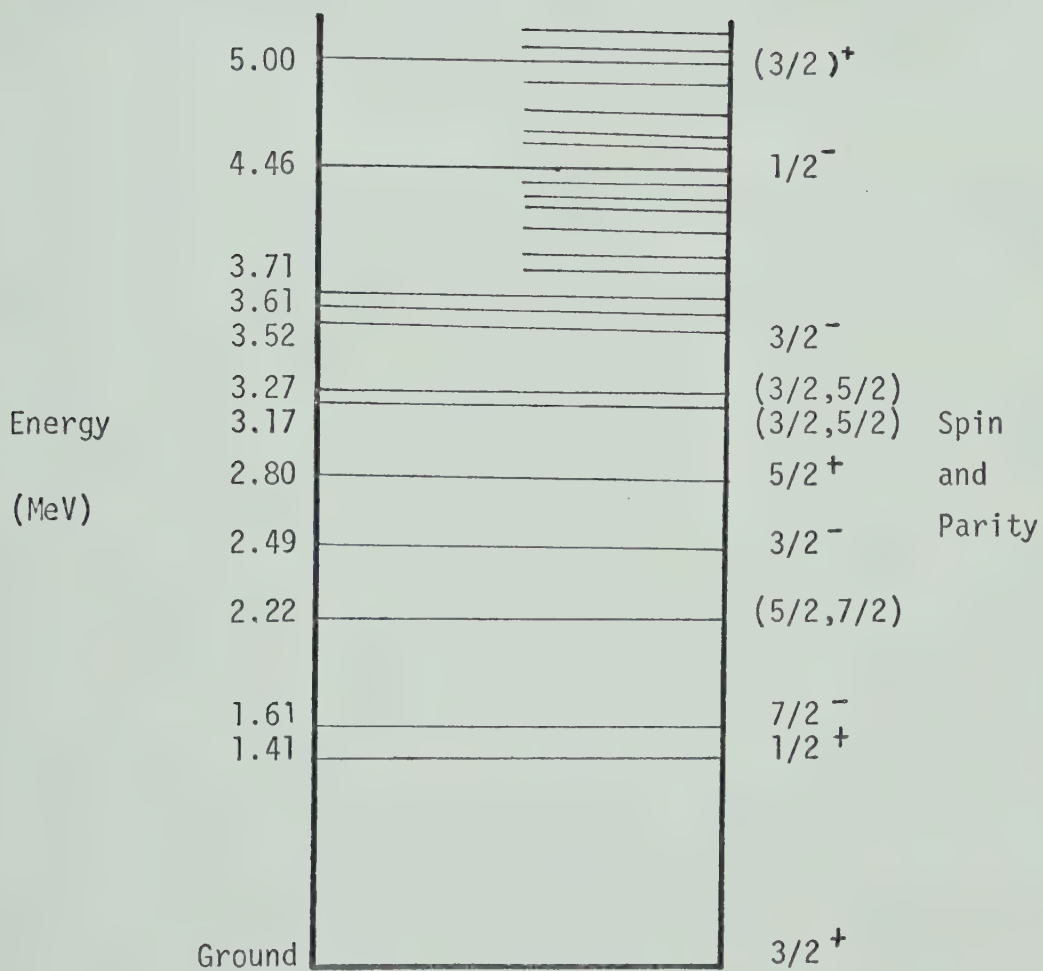
CHAPTER TWO

PREVIOUS STUDIES

2.1 Levels of Argon-37

The levels of Argon-37 are seen⁽¹²⁾ in figure 1. This nucleus can be regarded as a closely packed spheroid of 19 neutrons and 18 protons. The ground state configuration has been deduced⁽¹³⁾ from the ordering of levels in the shell model to be one neutron hole in the 1d 3/2 shell with a spin and parity of $3/2^+$. This ground state is unstable and decays to ^{37}Cl by electron capture with a half-life of 2.96×10^6 seconds⁽¹⁴⁾. On the other hand, the $3/2^+$ ground state of the mirror nucleus ^{37}K is also unstable, decaying to ^{37}Ar by positron emission. The "log ft" of this latter transition⁽¹⁵⁾ also leads to a $3/2^+$ value for the ground state of ^{37}Ar . This spin and parity has been further verified by studies⁽¹⁵⁾ of the atomic arc spectrum. This value gave a 43% spectroscopic factor to the observed d-wave neutron stripping in the $^{36}\text{Ar}(d,p)^{37}\text{Ar}$ reaction⁽¹⁵⁾. This spectroscopic factor has been verified by neutron pickup in the $^{38}\text{Ar}(d,t)^{37}\text{Ar}$ reaction⁽¹⁶⁾. Naudé et al have examined the normalized variance of cross section fluctuations in the $^{39}\text{K}(d,\alpha)^{37}\text{Ar}$ reaction. They also deduced⁽¹⁷⁾ a probable spin from this of $3/2$.

Several theoretical studies^(18,19,20) of the excited states of ^{37}Ar have been performed in the past but they are



^{37}Ar

FIGURE 1:
ENERGY LEVELS OF ^{37}Ar .

all somewhat unsatisfactory. The lowest level seems to be due to excitations of the unfilled s-d shell. Lande et al⁽²¹⁾ have attempted to describe the higher levels in terms of coupling between quadrupole oscillations of an even core and a neutron in nearly empty single particle levels.

The first excited state is at 1.41 MeV and has been seen with $^{37}\text{Cl}(p,n)^{37}\text{Ar}$ (22), $^{39}\text{K}(d,\alpha)^{37}\text{Ar}$ (23,24,25), $^{35}\text{Cl}(^3\text{He},p)^{37}\text{Ar}$ (26), $^{36}\text{Ar}(d,p)^{37}\text{Ar}$ (27,28,29,30), and $^{36}\text{Ar}(n,\gamma)^{37}\text{Ar}$ (32) reactions. This state has been populated⁽²⁷⁾ by s-wave neutron stripping by $^{36}\text{Ar}(d,p)$. Consequently it is a $1/2^+$ state. The spectroscopic factor is 13%, indicating that the $2s_{1/2}$ subshell was not completely full. This spin and parity was also deduced from the observed isotropy of the n_1 - γ angular correlation following the $^{37}\text{Cl}(p,n\gamma)$ reaction⁽³³⁾, and the spin from Naude's fluctuation analysis above⁽¹⁷⁾. Naude⁽³⁴⁾ has also compared experimental cross sections for $^{39}\text{K}(d,\alpha)$ to this level and others with theoretical Hauser-Feshbach cross sections normalized to the ground state. It was also consistent with a spin of $1/2$. The parity could not be uniquely determined by this method, however.

Since the Argon-37 nucleus has one hole in an otherwise full s-d neutron shell, one would expect some of the higher excited states to show a strong single particle component. This seems to be the case with the

second excited state seen at 1.61 MeV, with the same reactions listed for the previous level. The $^{36}\text{Ar}(d,p)$ reaction⁽¹⁵⁾ puts neutrons into the S.P. f 7/2 state at a spectroscopic factor of 82%. The f 5/2 value was eliminated because spin-orbit interaction places levels belonging to this subshell at 6.5 MeV⁽³⁸⁾ above the f 7/2 shell. A study⁽³⁵⁾ of the transition rate of γ decay of the 1.61 MeV level to the $3/2^+$ ground state verified the $7/2^-$ choice. It must be mentioned, however, that both Naude's fluctuation analysis⁽¹⁷⁾ and his Hauser-Feshbach fits⁽³⁴⁾ for the $^{39}\text{K}(d,\alpha)$ reaction favor 5/2 spin, but do not exclude 7/2 as a possibility. There seems to be some difficulty in applying certain aspects of the statistical theory such as the '2I+1 rule', to the higher spin states, especially if there are marked fluctuations and/or strong single particle components in the cross section^(24,36).

The 2.22 MeV level was observed in the $^{37}\text{Cl}(p,n)$ ⁽²²⁾, $^{39}\text{K}(d,\alpha)$ ^(23,24,25), $^{35}\text{Cl}(^3\text{He},p)$ ⁽²⁶⁾, and $^{36}\text{Ar}(n,\gamma)$ ⁽³²⁾ reactions. It was not located in the stripping spectra of $^{36}\text{Ar}(d,p)$ by either Rosner and Schneid⁽²⁷⁾ or Wiza et al⁽²⁸⁾, as a result no ℓ -value was assigned. However, γ rays seen by Howard et al⁽²⁹⁾ from this same reaction did testify to the presence of this level. Furthermore, their particle- γ ray angular correlations put its spin in the range (5/2, 7/2). Both Naude's fluctuation analysis⁽¹⁷⁾ and his Hauser fits⁽³⁴⁾ also limited the spin to the same values.

The 2.489 MeV level has been located in the same reactions listed under the first excited state, with the exception that it had not yet been seen in $^{37}\text{Cl}(p,n)$ until the present experiment. The $^{36}\text{Ar}(d,p)$ work of Rosner and Schneid⁽³⁷⁾, in particular, found it as one of a group of four excited states showing strong p-wave components. The 2.489, 3.52, 4.46, and 5.15 MeV levels have spectroscopic factors of 45, 36, 16, and 59% respectively. If the first two are assumed to be two forms of core coupling with the $2p\ 3/2$ S.P. level, and the last two, with the $2p\ 1/2$ S.P. level, then we can assign an average S.P. energy, weighted by the spectroscopic factors, to each subshell. The difference $2p\ 1/2 - 2p\ 3/2$ then becomes 2.09 MeV. This compares well with both the theoretical spin-orbit energy difference of 2.04 MeV⁽³⁸⁾ and similar energy differences in ^{41}Ar and ^{40}Ca ⁽³⁷⁾. A spin of $3/2$ for the 2.489 MeV level also results from the work of Naude above⁽¹⁷⁾.

The 2.80 MeV level was first seen⁽³⁹⁾ in the γ spectrum of ^{37}Ar following the positron decay of artificially excited ^{37}K . The study of the angular distribution of these γ rays to the ground state give a spin and parity of $5/2^+$. This level has also been seen in $^{39}\text{K}(d,\alpha)$ (23,24,25), $^{35}\text{Cl}(^3\text{He},p)$ (26), $^{36}\text{Ar}(d,p)$ (27,28,29,30), and $^{36}\text{Ar}(n,\gamma)$ ⁽³²⁾ work. The $5/2$ spin was confirmed by Naude's work⁽¹⁷⁾ also.

Two close levels at excitation energies of 3.17 and 3.27 MeV were observed with $^{39}\text{K}(\text{d},\alpha)^{(23,25)}$, $^{35}\text{Cl}(^3\text{He},\text{p})^{(26)}$, and $^{36}\text{Ar}(\text{n},\gamma)^{(32)}$ reactions. Also, the particle- γ angular correlations from $^{36}\text{Ar}(\text{d},\text{p}\gamma)^{(29)}$ limited the spin to (3/2, 5/2) for both levels. The observations of a 0.70 MeV level by $^{34}\text{S}(\alpha,\text{n})^{(40)}$ and a 2.41 MeV level by $^{37}\text{Cl}(\text{p},\text{n})^{(22)}$ have not been duplicated in any later work, either with identical reactions or any other reactions. They are consequently not listed in recent tables⁽¹²⁾ of Argon-37 levels. The density of known levels further increases at higher energies, but their discussion is beyond the scope of this work.

The ^{37}K nucleus is the mirror nucleus of ^{37}Ar . On the basis of energy, spin, parity and sometimes similarity of γ decays, the states 0, 1.37, 1.38, 2.28, 2.17, and 2.75 MeV of ^{37}K seem to be mirrors of the states 0, 1.41, 1.61, 2.22, 2.489 and 2.80 MeV of ^{37}Ar respectively⁽⁴¹⁾. If we assume these mirror nuclei differ only in Coulomb energy, an average Coulomb energy difference of 6.09 MeV follows from these energy levels. Since the proton distribution is not so spherical in the odd proton ^{37}K as in the odd neutron ^{37}Ar , this difference is smaller than the theoretical value of 8.4 MeV if both nuclei were spherical⁽⁴²⁾.

2.2 $^{37}\text{Cl}(p,n)^{37}\text{Ar}$ Reaction

The $^{37}\text{Cl}(p,n)$ reaction was chosen because work remains to be done with it on those levels below the 2.489 MeV level and no previous work has been done at or above this level. The proton beam energy was around 5 MeV as compared with a Coulomb barrier of 5.9 MeV for protons on Chlorine-37. It was expected that we could compare the cross sections to statistical model predictions with little or no direct interaction contribution. Previous work⁽³³⁾ with this reaction up to 5.5 MeV proton energy, did not show any direct interaction effects.

The resonances observed⁽⁴³⁾ in $^{34}\text{S}(\alpha,\gamma)^{38}\text{Ar}$ at around 10 MeV excitation energy put an upper limit on the average Argon-38 compound nucleus level separation, D , equal to 60 keV at that energy. This was too large, owing to the fact that α and ^{34}S can form only levels of parity $(-1)^J$. The $^{37}\text{Cl}(p,\gamma)^{38}\text{Ar}$ ⁽⁴³⁾ reaction showed 196 resonances in the proton energy range, 0.427 - 1.846 MeV, which gave an upper estimate to D of 4.8 keV at $E_x = 12$ MeV. The $^{37}\text{Cl}(p,n)^{37}\text{Ar}$ reaction just above threshold⁽⁴⁴⁾ had resonances with an upper limit to D of 8.1 keV for $E_x = 12$ to 13 MeV. The density of compound nucleus states should increase exponentially with the square root of the excitation energy if the nucleus behaves as a Fermi gas⁽⁴⁵⁾.

The level width, Γ , in Argon-38 also increases rapidly with energy as more and more channels become

available to decay. The α , p, and n channels open up at 7.2, 10.2, and 11.8 MeV excitation, respectively. The width broadens still more as the transmission functions for the channels through the potential barriers approach unity. Consequently, one would expect the level widths to approach D a few MeV above threshold. When these levels begin to overlap, it is no longer possible to directly observe them in the excitation function regardless of resolution, instead, we have fluctuations caused by the interference of the overlapping levels. There is then no correlation between the energy variation of total cross sections at neighboring energies nor of the cross sections to each channel at the same energy. In fact, the symmetry of decay products about 90° to the beam, exhibited by all single compound nuclear resonances undergoing decay, is destroyed by interference between overlapping levels of different parities.

Von Witsch et al⁽⁴⁶⁾ observed these effects in the Argon-38 nucleus when it was excited to around 21 MeV with the $^{37}\text{Cl}(p,\alpha)^{34}\text{S}$ reaction. They measured the coherence energy of the fluctuations to be 13 keV at $E_x = 20.2$ MeV and 18 keV at $E_x = 21.0$ MeV. This is a measure of the average total width, Γ , of the interfering levels assuming no spin or parity dependence.

The symmetry in the angular distributions is restored if the cross sections are averaged over an

energy interval much larger than the coherence energy of these fluctuations. This is because the interference terms will then average out to zero (see the Random Phase Approx. in section 1.2). By comparing this distribution with Hauser-Feshbach theory, using independent calculations for the transmission coefficients for the incident and emergent channels, one can arrive at the sum of the transmission coefficients over all the mostly unknown channels (see equation 16). The equation 32, is then used, plus the assumption of spin distribution of:

$$(21) \quad D^J = D^0 / (2J+1) \exp(-J(J+1)/2\sigma^2)$$

to get the ratio Γ/D^0 for Argon-38.* D^J is the spacing of levels of spin J and σ^2 is the spin cut-off parameter. The effect of σ^2 is mainly felt in the shape of the angular distribution. Von Witsch et al⁽⁴⁷⁾ deduced for ^{38}Ar at $E_x = 21$ MeV, that $\sigma^2 = 10^{+5}_{-2.5}$ and that $\Gamma/D^0 = 27$. Consequently $\Gamma/D \cong 600$ for the above distribution at the specified value for σ^2 . Combining this value with their Γ above, resulted in a $D = 30$ ev at $E_x = 21$ MeV.

The level density as a function of energy is fairly well known in Argon-38 up to 6.9 MeV from counting the individual levels⁽⁴⁸⁾. This function was extrapolated to pass through the point corresponding to Von Witsch's value of $D = 30$ ev at 21 MeV, to get the level density at an

* Γ is assumed to be independent of J .

excitation energy of 15.2 MeV corresponding to $E_p = 5$ MeV. It can be shown from thermodynamics⁽⁴⁵⁾ that for a Fermi gas, the density of states:

$$(22) \quad \rho(E_x) \cong \exp 2 \sqrt{aE_x} \quad ,$$

where a is a constant. This means that a plot of $\ln \rho(E_x)$ v/s $\sqrt{E_x}$ will be a straight line of slope $2\sqrt{a}$. Using this, an estimate of $\rho(15.2 \text{ MeV})$ was arrived at, viz. 3,400 levels/MeV or $D \cong 0.3 \text{ keV}$, and with an averaging interval of 160 keV, for instance, the cross sections will be averaged over 500 resonances. Thus the statistical theory was expected to be valid at this energy.

The slope of such a plot gave $a = 3.4 \text{ MeV}^{-1}$. This compared well with Cameron's formula⁽⁴⁹⁾ for $a = 0.095 \times$ (the atomic mass), which gave in this case, $a = 3.6 \text{ MeV}^{-1}$. Von Witsch et al⁽⁴⁷⁾ obtained $a = 4.6^{+0.2}_{-0.1} \text{ MeV}^{-1}$. They, however, used a more complex level density function due to Lang and Le Couteur⁽⁵⁰⁾, so the two results are not comparable.

2.3 Excitation Function

Barnard et al⁽⁵¹⁾ observed 'resonances' in the $^{37}\text{Cl}(p,n)^{37}\text{Ar}$ reaction from threshold to $E_p = 4 \text{ MeV}$. Two such features at $E_p = 2.1$ and 2.3 MeV and a few hundred keV wide appeared under higher resolution ($\Delta E = 1.7 \text{ keV}$

in this case) to be the envelopes of hundreds of individual compound nuclear states⁽⁴⁴⁾. They postulated that these are due to the existence of a giant resonance, a few keV above threshold with an anomalous small width which had been split by spin-orbit coupling. They suggest the small width arose from the very small imaginary potential for nearly doubly magic nuclei. (See also Lane, et al⁽⁵²⁾).

Barnard et al⁽⁵¹⁾ observed similar features in the excitation function of neutrons to the ground state of ^{37}Ar at $E_p = 2.9, 3.3, 3.4, 3.6$, and 3.75 MeV and at $E_p = 3.7$ and 4 MeV in the excitation function of neutrons to the first excited state of ^{37}Ar . Thus the observed structure in the excitation function appears to be channel dependent. Iyengar et al⁽³³⁾ measured the excitation function of γ -rays from the first excited state ^{37}Ar to its ground state following the $^{37}\text{Cl}(p,n)$ reaction. Their target was 150 keV thick to 5 MeV protons. They too observed a resonance at $E_p = 3.6$ MeV and a possible rise in the yield at 4.0 MeV. In addition, sharp peaks were seen by them at $4.5, 5.1$, and 5.4 MeV. Furthermore, their neutron angular distributions at $E_p = 5.1, 5.3$ and 5.5 MeV from targets 100 keV thick to a 5 MeV proton beam do not show symmetry about 90° . The random phase approximation seems to be not fulfilled in this energy range in view of those resonances⁽⁵³⁾. As we shall see, further reservations on the use of the statistical theory are also concluded from the present work.

CHAPTER THREE

EXPERIMENTAL DETAILS

3.1 Beam Injection and Acceleration

The $^{37}\text{Cl}(p,n)^{37}\text{Ar}$ reaction was studied with the time-of-flight technique, using the pulsed beam provided by the University of Alberta Van de Graaff accelerator facility. The parts of the accelerator are presented in figure 2. The accelerator, built by the High Voltage Engineering Corporation, is an electrostatic generator where the positive charge is mechanically transported by a conveyor belt to the high voltage terminal, resulting in a self-focusing field with up to 6.5 megavolts potential drop along the acceleration tube.

This tube was fed with a low energy proton beam produced by the D.C. field set up between the probe electrode and extraction canal of the source bottle. The field was selected to give maximum emission from this source, and optimum conditions for focusing with the electrostatic lenses (Einzel lens and Focus 3) and the acceleration tube. The ion type was selected and aimed by the terminal analysis magnet after the appropriate gas had been injected into the source. This was ionized in the low pressure gas discharge produced in the 125 MHz R.F. field maintained across the source. The source

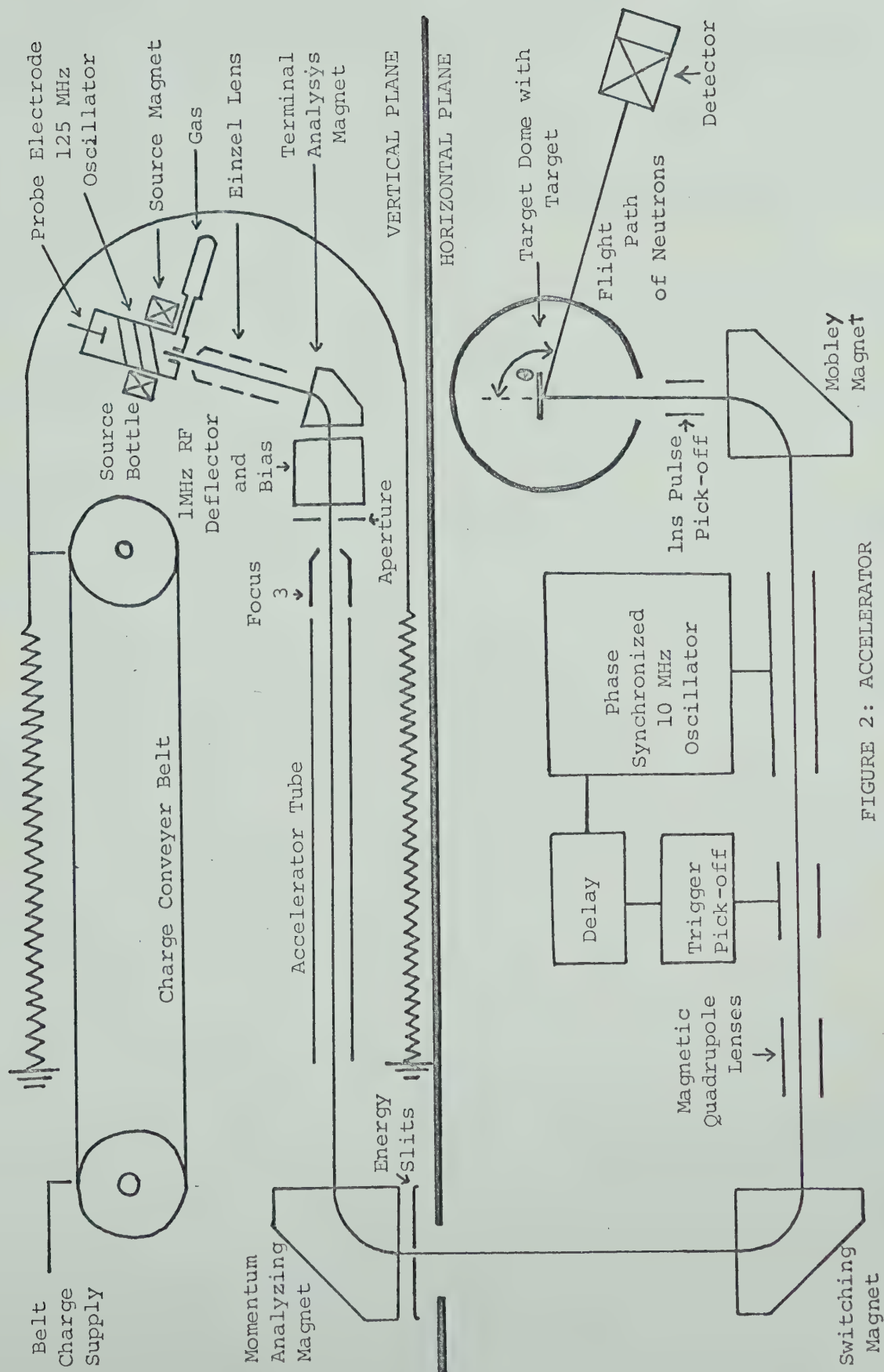


FIGURE 2: ACCELERATOR

magnet intensified this discharge. Its axial magnetic field gave a helical motion to the discharge electrons on their way to the probe electrode, increasing their path through the gas. The source can provide a 2 milliampere D.C. proton beam.

The beam was pulsed in the following manner. It passed through two sets of R.F. deflection plates at right angles to each other, which were driven 90° out of phase with a frequency of 1 megahertz. These plates were D.C. biased to permit the resulting elliptical beam path to cross a 1/16" aperture. The resulting high intensity bursts of 10 nanoseconds length were then compressed for nanosecond time-of-flight work by post-acceleration bunching. In this way, a 1 megahertz pulsed proton beam of 10 microamperes average current could be injected into the acceleration tube. The above equipment was electrically isolated in the field-free region of the dome-shaped high voltage terminal and powered from the conveyor belt pulley.

3.2 Beam Transport System

The accelerated protons were transported to the target facility through the evacuated drift tube, successively aimed and focused by the momentum analyzing magnet, the switching magnet, the magnetic quadrupole lenses, and the Mobley magnet (see figure 2). Concrete, water,

paraffin, and/or lead structures shielded workers from any neutrons and γ rays produced from interactions of the accelerator beam with the sides of the drift tube and the beam defining apertures at various stages along its path.

The beam energy stability was maintained by corona current from the high voltage terminal which was automatically controlled to reduce the difference between beam current striking the two energy slits. These slits are located at the focus of the 90° momentum analyzing magnet and together with it maintained the energy resolution of the beam. The strength of the analyzing magnet's field, and hence the energy of the protons passing through the slits, was calibrated to the frequency of nuclear magnetic resonance of a container of hydrogen gas inserted into the field. This energy calibration was based on observations of the threshold of the ${}^7\text{Li}(p,n){}^7\text{Be}$ reaction. It was reliable to ± 3 keV at 6 MeV beam energy.

The 10 nanosecond pulses were compressed by the Mobley bunching system. It contains a 10 megahertz oscillator which can produce a R.F. deflection field between plates located at the first focus of the 90° Mobley magnet. This oscillator was synchronized to the pulses by the trigger signal from the cylindrical pick-off upstream. The field swept each pulse such that the leading and lagging particles of the pulse were deflected into

longer and shorter paths, respectively, through the Mobley magnet. These paths had a mean radius of curvature equal to 75" and converged at the second focus where the target was located. This combination of pre-acceleration pulsing and post-acceleration bunching can produce 3 microamperes average current of pulsed beam at the target of repetition frequency of 1 MHz, each pulse being of 400 picoseconds duration^(54,55).

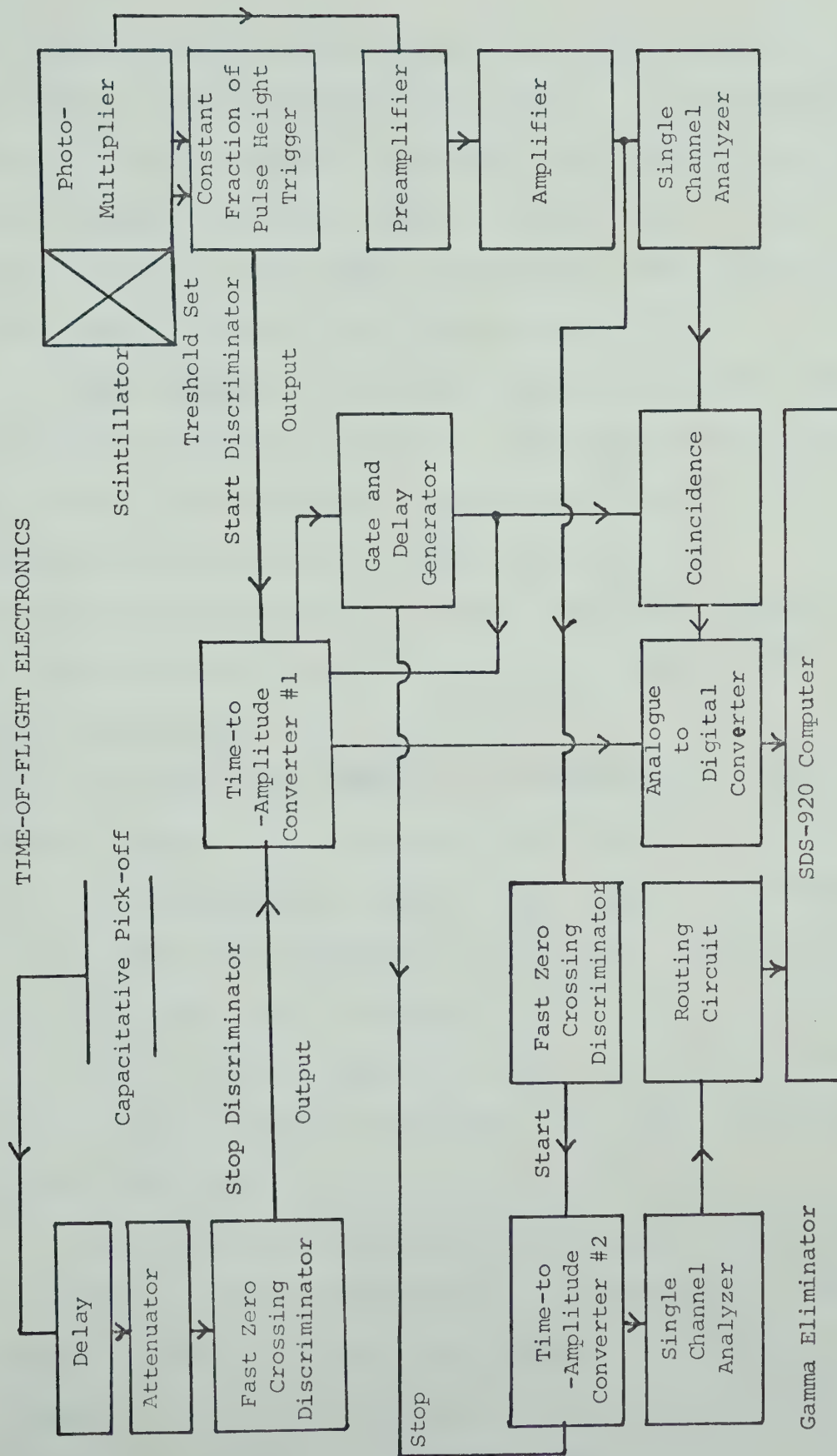
3.3 Neutron Spectrometer

The target was located in the center of an evacuated steel dome. The neutron energy groups emitted from the target were seen by electronically measuring their flight time over an accurately known distance ('d'). This method can provide for good time resolution for such neutron energies as 100 keV to 20 MeV. The neutron detector was located on a cart whose distance from the target could be remotely controlled along a boom over a range of 1.0 - 6.3 meters. The angle of the boom to the beam line could also be remotely varied continuously over the range 0° - 150° .

The timing was started when the start discriminator signalled the passage of a neutron through the detector (see figure 3). 30 cm ahead of the target is a 3 cm cylindrical capacitor. The stop discriminator detected

Figure 3

TIME-OF-FLIGHT ELECTRONICS



the passage through this pick-off of that beam pulse which followed the one producing the neutron. This stopped the timing. This arrangement reduced electronics' dead time since timing began only after a neutron event in the detector and not with every beam pulse. The jitter due to the Mobley system was ≤ 40 picoseconds⁽⁵⁶⁾.

The timing was done by an ORTEC Model 437A time-to-amplitude converter (T.A.C.-1). This conversion was linear. The T.A.C. output was analyzed by a TMC Model 217A analogue to digital converter with 4,096 channels (ADC-A). This TAC-ADC-A combination can provide long term stability and negligible shifts in peak position for high counting rates while maintaining a very low noise level over the 1000 nanosecond time spectrum⁽⁵⁵⁾. It was coupled to an on-line SDS-920 computer which operated on a general purpose collection and analysis code (GPKS). This code's subroutines allowed continual displaying of the time spectrum during data collection, the recording of this data on magnetic tape, and such preliminary analyses as the determination of peak energy, total counts under the peak, subtraction of background, etc.

The combined conversion from time to channels was calculated in the following way. A 10 megahertz time mark generator, feeding a fast-zero crossing discriminator, produced a train of sharp start pulses for the T.A.C. Stop pulses, bearing an integral relation to the start

pulses, were formed by a second start pulse output from the above discriminator in coincidence with lower frequency pulses. These latter could be either the output of a second generator or random pulses from neutron events in the detector. The resulting time spectrum consisted of a series of peaks, 100 nanoseconds apart, whose separation in channels was easily calculated from the 920 display. The T.O.F. time spectrum was then calibrated by relating this number of nanoseconds per channel to the observed gamma peak, whose flight time, d/c , was known.

3.4 The Neutron Detector

The neutrons were detected by the fluorescence produced in a liquid organic scintillator, type NE-218. The fluorescence was produced indirectly, through their elastic collisions with hydrogen molecules which, in turn, interacted with the fluor in the liquid. The result was a light pulse (of several nanoseconds duration e.g. 5.6 nsec. for the 'fast component' ⁽⁵⁴⁾), detected using a photomultiplier tube. The light intensity emitted was a smooth but nonlinear function of the energy imparted to the protons. The proton energy, however, could be any fraction of the neutron energy, depending on the angle of the recoiling proton with respect to the neutron, thus the necessity of timing the pulse's leading edge. However, the neutron transit time through the 3/4 inch

thick scintillator, which is 700 picoseconds for a 4 MeV neutron (and 1,400 picoseconds for a 1 MeV neutron), was a factor in limiting time resolution. The photomultiplier (P.M.) could have 300 picoseconds of intrinsic time uncertainty, depending on the individual tube and the potentials of its focusing grids. This is related to both the statistics of electron emission from the photocathode and variations in their transit time from different parts of the cathode to the first dynode. To minimize the time dispersion due to the entire detector assembly, a finite threshold was set for the start discriminator, such that P.M. output pulses below that level were not accepted.

To obtain absolute cross sections, the efficiency of the scintillator-photomultiplier-discriminator combination had to be determined. The amplitude spectrum of P.M. output pulses for a given neutron energy could be theoretically determined. This distribution would be a combination of the proton recoil spectrum and the emission statistics of the photocathode. The efficiency is then just the integral of this distribution above the threshold. These calculations were performed using the program, Effi⁽⁵⁷⁾, at the University of Alberta.

The neutron detector threshold was determined by observation of the amplitude spectrum from a gamma source

such as ^{22}Na . Each gamma peak also produces a broad range of light intensities above the threshold as a result of Compton collisions with the material of the scintillator. The energy for threshold electrons was then determined from the known energies of the 'Compton edges' and later related to the proton energy corresponding to the threshold, through a graph⁽⁵⁸⁾ of proton/electron response for the scintillator.

3.5 Timing Electronics

The general aspects of the time-of-flight technique were outlined in Section 3.3. Figure 3 shows the details of this system. The stop discriminator, a fast zero crossing discriminator (FZCD), produced a timing pulse at the cross-over point of the beam pick-off current. The walk or variation in timing was ± 65 picoseconds over a 30:1 dynamic range of signal amplitude⁽⁵⁶⁾. The start discriminator, a constant fraction of pulse height trigger (CFPHT), produced a timing signal at the leading edge of the P.M. anode pulse, which was based on the time this pulse reaches a constant fraction of its maximum value. This fraction was chosen to be 10% to minimize the time resolution for almost all output amplitudes and reduce the walk to ~ 50 picoseconds over an amplitude range corresponding to 500 keV - 3 MeV neutrons⁽⁵⁹⁾. The CFPHT

utilized an attenuation-subtraction technique which resulted in the desired zero-crossing pulse. This, in turn, fed a FZCD which produced the timing signal. The FZCD was effectively gated with the pulse from the last P.M. dynode, whose amplitude had to exceed the bias threshold set on a tunnel diode in the circuit for a signal to be sent.

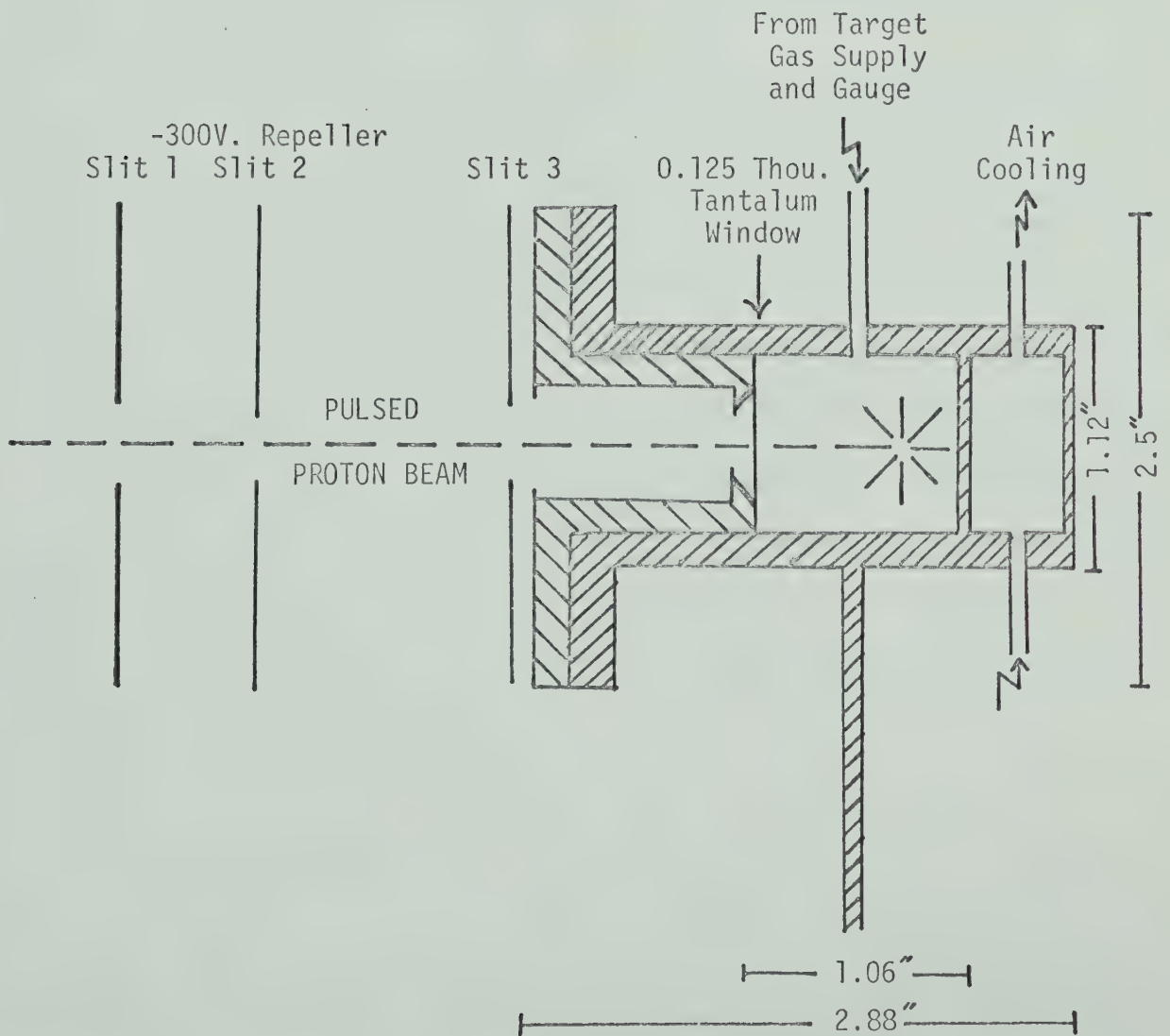
The background to the time spectrum can be reduced by a gamma elimination system⁽⁵⁵⁾. This circuit was based on the difference between γ pulse shapes and that of neutrons. The γ pulses had longer tails. A measure of this shape was the time, between the arrival of the start discriminator signal from the leading edge of a pulse, and the FZCD signal from the cross-over point of the corresponding differentiated output of one of the P.M. dynodes. This time difference was 10 - 20 nanoseconds. The time was measured and selected by the T.A.C.-2 and single channel analyzer combination of figure 3. If a neutron was detected, a routing circuit instructed the on-line computer to store it in a separate region of computer memory. Any neutrons incorrectly routed were easily seen in the original spectrum of the T.A.C.-1 - ADC-1 combination. It was important that the output of the two T.A.C.'s, as well as that of the amplitude gate, seen in figure 2, be synchronized. The gate and delay generator did this by timing the output of the T.A.C.-1 to the output of the CFPHT.

3.6 Chlorine-37 Targets

The reaction was investigated using several different targets. Solid target #1 consisted of NaCl, evaporated in vacuum onto a tungsten foil. The NaCl was $90.36 \pm 0.10\%$ artificially enriched in the ^{37}Cl isotope. Solid target #2 was similar to the above but on a tantalum foil. These were thin targets, the latter being 17 keV thick to a 5 MeV proton beam, and were used to study neutron angular distributions. Most of the evaporated targets were poorly bonded to their backings and pieces flaked off under a 10 watt beam, making absolute cross section measurements impossible.

Consequently a gas target was constructed to measure these cross sections. It had the added advantage that its thickness could be easily varied from lower thickness values required for obtaining excitation functions, to the higher thicknesses needed for obtaining angular distributions averaged over a wide energy range. Natural ethyl chloride was successfully used as the gas (Natural chlorine is 24.4% ^{37}Cl). The beam entered the gas cell through a $\frac{1}{2}$ " wide tantalum window which is 125 microinches thick, or 150 keV to a 5.45 MeV proton beam. (See figure 4). The 1.063" deep cell contained ethyl chloride at a pressure of 25.5 cm Hg (or 160 keV thick to the 5.3 MeV beam emerging through the window) used in the angular distributions. The temperature of the air cooling was monitored and found

FIGURE 4:
GAS TARGET



to be 23°C . The cell should be kept sealed during the experiment to prevent the density of the gas from varying due to heating of the gas by beam bombardment. The cell was tested for leakage by measuring the gas pressure after the beam was shut off and the cell had cooled to room temperature.

The beam was collimated onto the axis of the cell by three slits positioned as shown in figure 4. The diameters of the circular slits were $1/4"$, $3/8"$, and $3/8"$ wide at distances of 12 cm, 3.6 cm, and 0.3 cm, respectively, from the front of the cell. In this way, it was ensured that the beam incident on the cell entered through the window. The current incident on the cell was measured with a current integrator which was the beam current passing through the gas. Slit 2 was biased to -300 volts to repel secondary electrons coming from the beam defining slit #1.

3.7 Energy Resolution

The resolution was affected by the target and the cell containing it. The gas target contributed an uncertainty in energy equal to the energy loss of protons passing through it, i.e. 160 keV in the above example. The cell itself contributed to energy spread in two ways. The energy straggling in the window had a standard devia-

tion of 18 keV for a 5 MeV proton beam⁽⁶⁰⁾. The second factor was due to the finite length of the cell, L , and is given by:

$$(23) \quad \Delta t \approx 72.3 L \left(\frac{1}{\sqrt{E_m}} - \frac{\cos \theta}{\sqrt{E_n}} \right), \quad \text{in nsec.}$$

L is in meters, E_m is the mean beam energy in the cell, and E_n the mean neutron energy, both in MeV, and θ is the reaction angle. The angular dependence is an important feature, with Δt varying from -0.12 to $+1.75$ nanoseconds for $\theta = 0^\circ$ to 150° when $E_m = 5.22$ MeV and $E_n = 4$ MeV. The algebraic signs are due to the coherence between this error and the gas target thickness error so that the total error is actually reduced at forward angles. Even with a target as thick as 160 keV, this error must be included at back angles with higher neutron energies. The contributions from beam width, window straggling, scintillator transit time and the intrinsic photomultiplier error, on the other hand, were small in comparison (see Sections 3.2 and 3.4). For $E_n = 4$ MeV the total uncertainty in energy ranged from 150 to 260 keV with angles 0° to 150° . At $E_n = 1$ MeV, on the other hand, this uncertainty was independent of angle and completely determined by the target thickness of 160 keV.

3.8 Cross Sections

The observed number of neutrons above the background was corrected for instrument dead time and detector efficiency to get the true number of neutrons passing through the scintillator. The differential cross section was easily calculated from the formula:

$$(24) \quad \frac{d\sigma}{d\Omega} \text{ (in mb)} = 3.39 \times 10^{-10} \times \frac{R^2}{d^2} \times \frac{\text{atomic weight of target}}{\text{no.gms/cm}^2 \text{ of target}} \\ \times \frac{N}{Q} \times \frac{1}{E} \quad .$$

R is the detector distance; d, the scintillator diameter; N is the number of neutron counts corrected for dead time; Q, the total charge due to beam incident on the gas* and E is detector efficiency. The transformation was then made to center-of-mass coordinates.

To allow for any variations in target thickness or beam direction, the solid targets were normalized to a monitor detector, placed at a fixed distance and angle from the target. This monitor fed a simpler version of the time-of-flight system than that described for the main detector (see Section 3.5). The angular distributions from the gas cell, however, were not so normalized, since the monitor used did not have a fixed efficiency. The instability was subsequently traced to its constant fraction of pulse height trigger. So the whole of equation

* Q is in μcoul .

24 was used, based on the integrated beam current reaching the target. For this to be done, the beam must be well defined by the three slits and the target stability ensured by its constant gas pressure. The inadequacy in the former working assumption is discussed in Chapter 4.

CHAPTER FOUR

EXPERIMENTAL RESULTS

4.1 Solid Target Results

The levels of Argon-37 discussed in Chapter 2 can be seen in the spectrum of figure 5. This spectrum is from a NaCl target (target #1 of Section 3.6) placed at a 0° angle to a beam of 5.75 MeV protons. The 3.27 MeV level has been recently isolated⁽⁶¹⁾ by using this reaction with a PbCl_2 target. Levels above 2.22 MeV in Argon-37 are seen here for the first time via the $^{37}\text{Cl}(p,n)$ macrochannel. As can be seen, there is no evidence here of the 2.41 MeV level that Ferguson and Paul⁽²²⁾ have claimed with the same reaction. Bernard et al⁽⁵¹⁾, also with this reaction, have observed the intensity of the 1.61 MeV level of ^{37}Ar as being $\leq 15\%$ of the group to the 1.41 MeV level. However, our figure shows a comparable intensity. In fact, at the lower mean proton energy of 5.22 MeV, discussed below, the ratio of the cross sections of the 1.61 to the 1.41 MeV levels are very nearly equal to the ratio of their respective values of $2I+1$, where I is the spin of the level.

In figures 6-8 are the angular distributions which were taken at mean proton energies of 4.95 MeV, 5.00 MeV and 5.05 MeV respectively for the first three levels of

Intensity (Relative Units) ↑

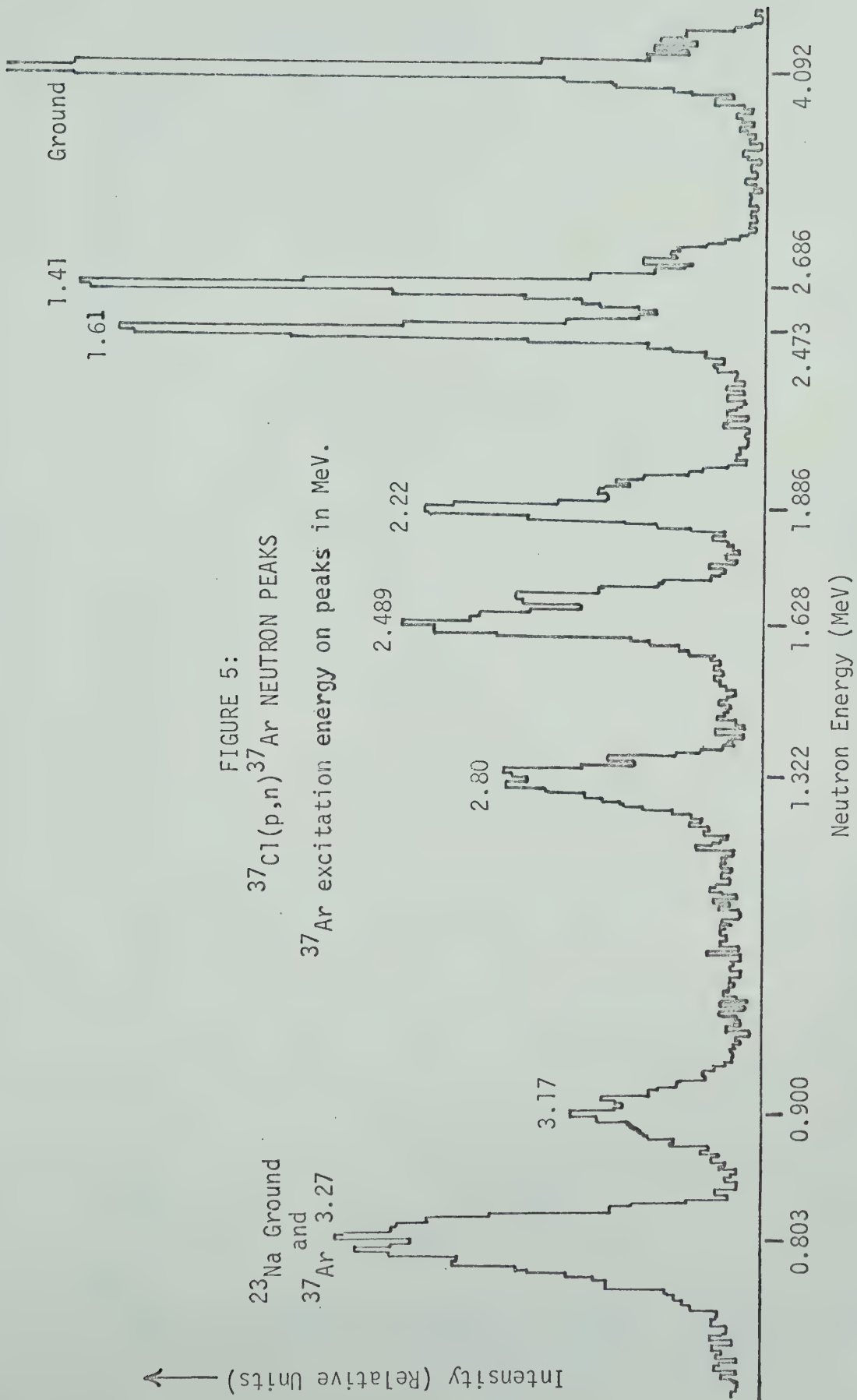


FIGURE 6: $^{37}\text{Cl}(p,n)^{37}\text{Ar}$ GROUND STATE NEUTRONS'
ANGULAR DISTRIBUTIONS AT THREE ENERGIES

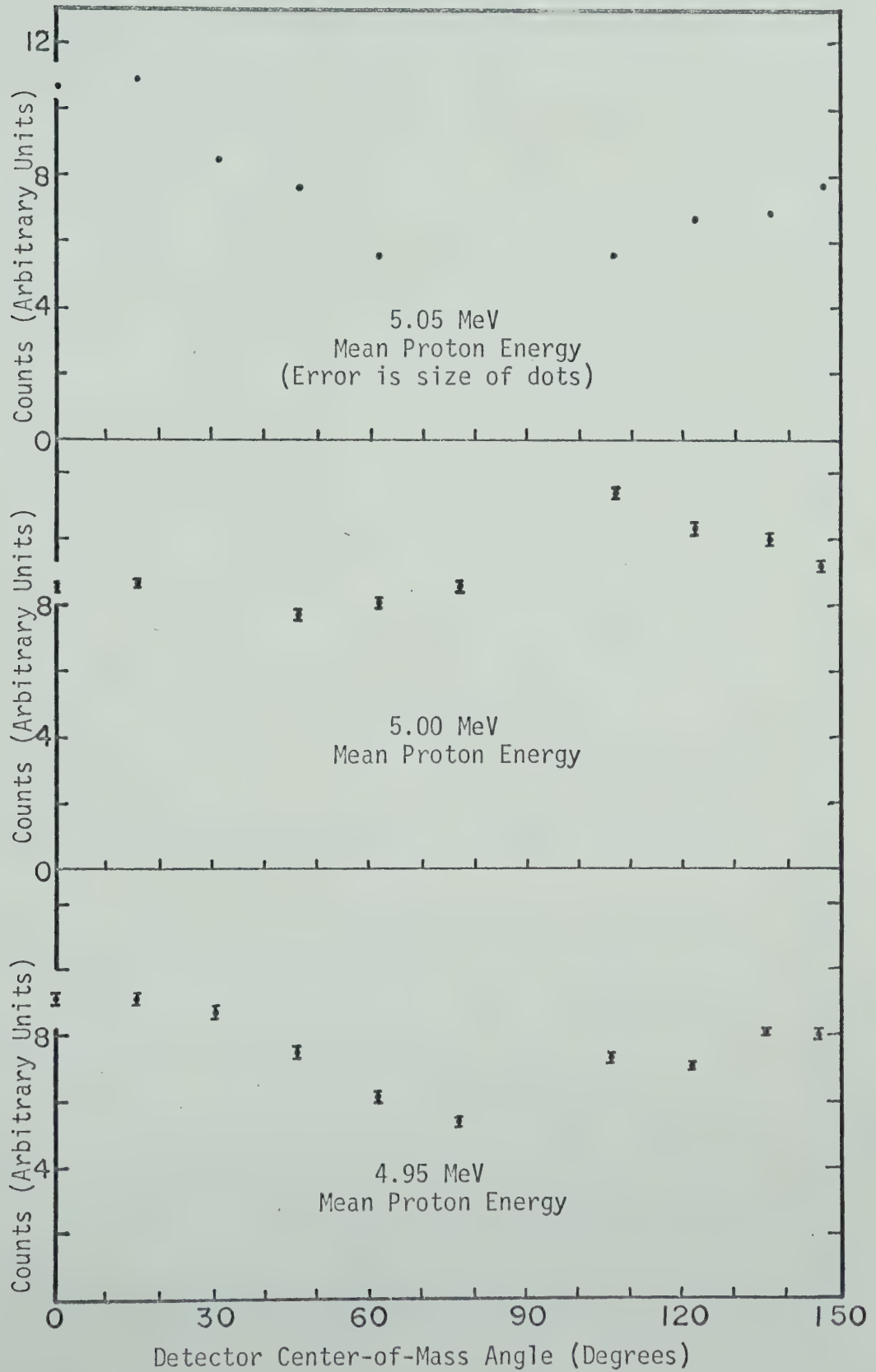
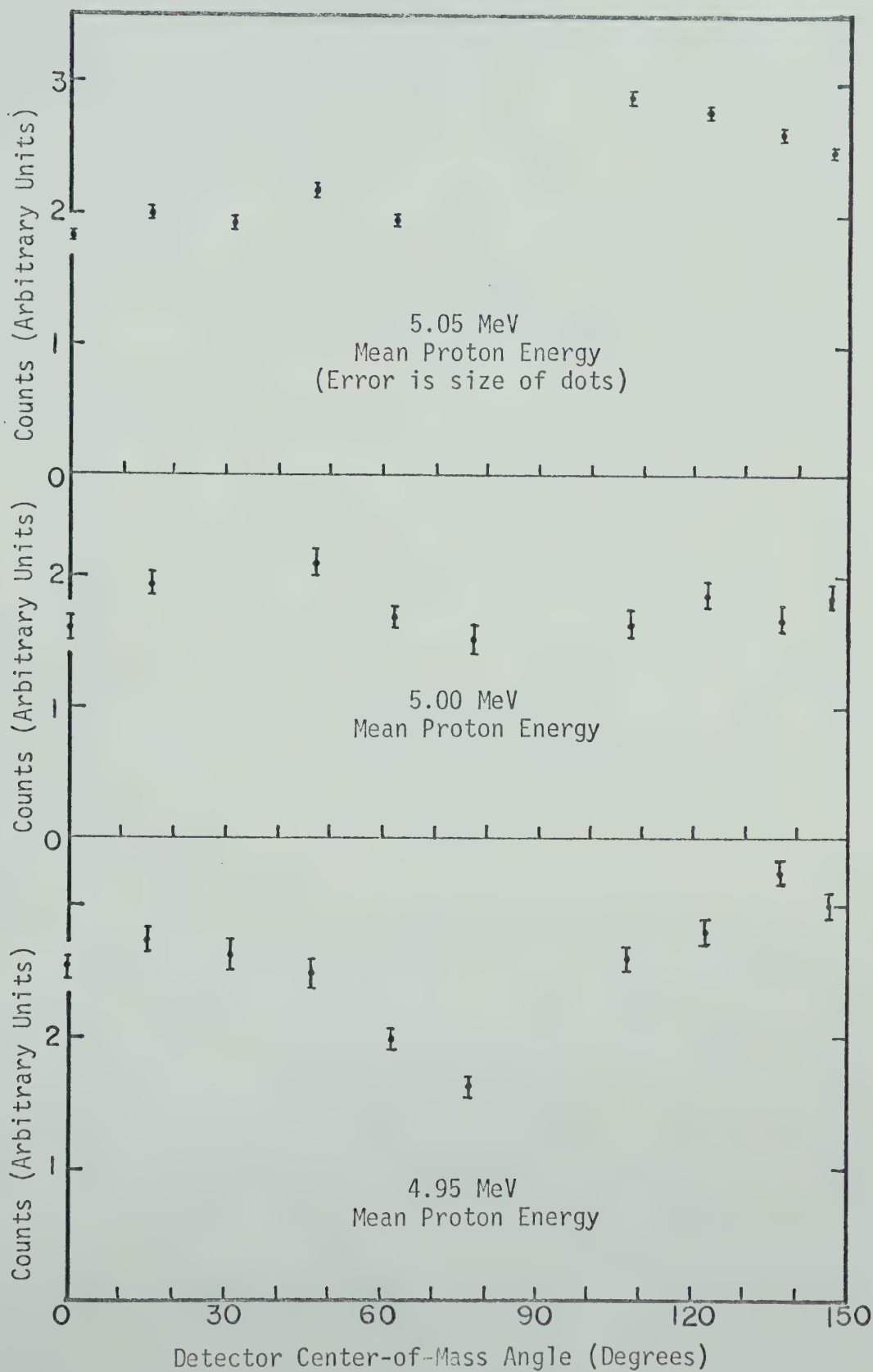


FIGURE 7: $^{37}\text{Cl}(p,n)^{37}\text{Ar}$ 1.41 MeV EXCITED STATE
NEUTRONS' ANGULAR DISTRIBUTIONS AT THREE ENERGIES



³⁷Ar. These distributions were normalized with respect to a detector fixed at 30° to the beam. The #2 target was 17 keV thick to a 5 MeV proton beam and so is thin enough to show asymmetry about 90° due to resonances of this order of size. Also seen are the sharp variations of the shape of the neutron angular distributions with bombarding energy. A resonance at 5 MeV mean proton energy had been located in the γ -ray yield by Iyengar et al⁽³³⁾ in the same reaction (see Section 2.3).

4.2 Gas Target Results

Ethyl chloride, 160 keV thick to the beam was used as a target in the gas cells. The mean proton energy in the cell was 5.22 MeV. As was discussed in Section 3.8, the monitor detector, to which the resulting neutron angular distributions were to be normalized, failed to have a fixed efficiency. The cross sections were thus calculated on the basis of the integrated beam current reaching the target. This can only be done, however, if the beam is well defined by the slits. Otherwise random variations in focusing conditions in the beam transport system would result in fluctuating amounts of charge striking the sides of the cell assembly and being recorded as part of the beam current, even though this charge had not passed through the gaseous target. Two detectors were

used to collect data and by the experimental set-up were exposed to the same angle at different runs. The resulting large variations in computed cross section observed between the two detectors can only be explained by this charge error. Re-runs of some of the angles confirmed the poor reproducibility of the results.

This charge error was reduced by taking the average result from the detectors at a given angle, this average being weighted inversely to the variance of these results.

$$(25) \quad \bar{x} = \frac{\sum_i \frac{1}{\epsilon_i^2} x_i}{\sum_j \frac{1}{\epsilon_j^2}} .$$

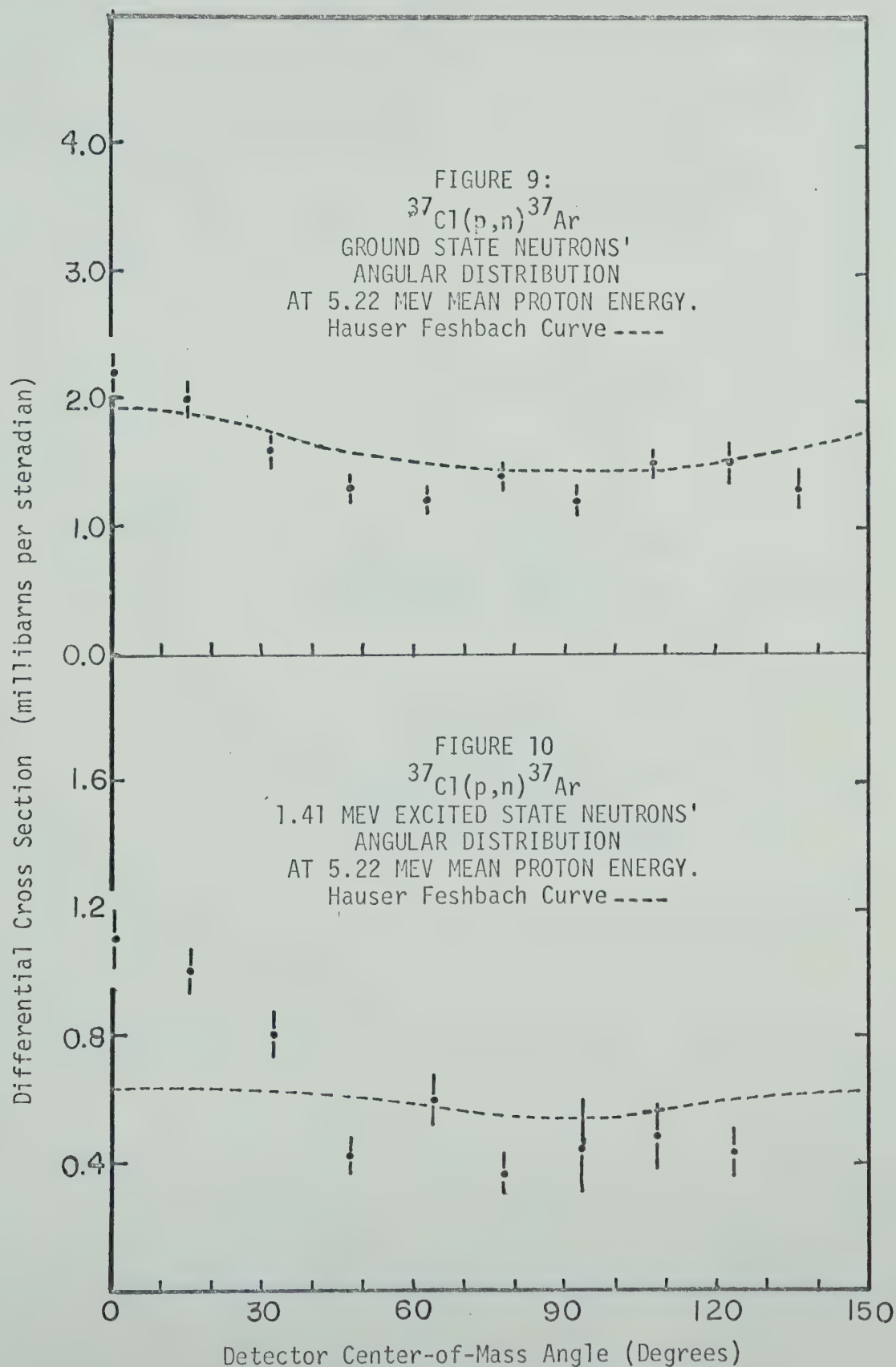
The error of the mean is

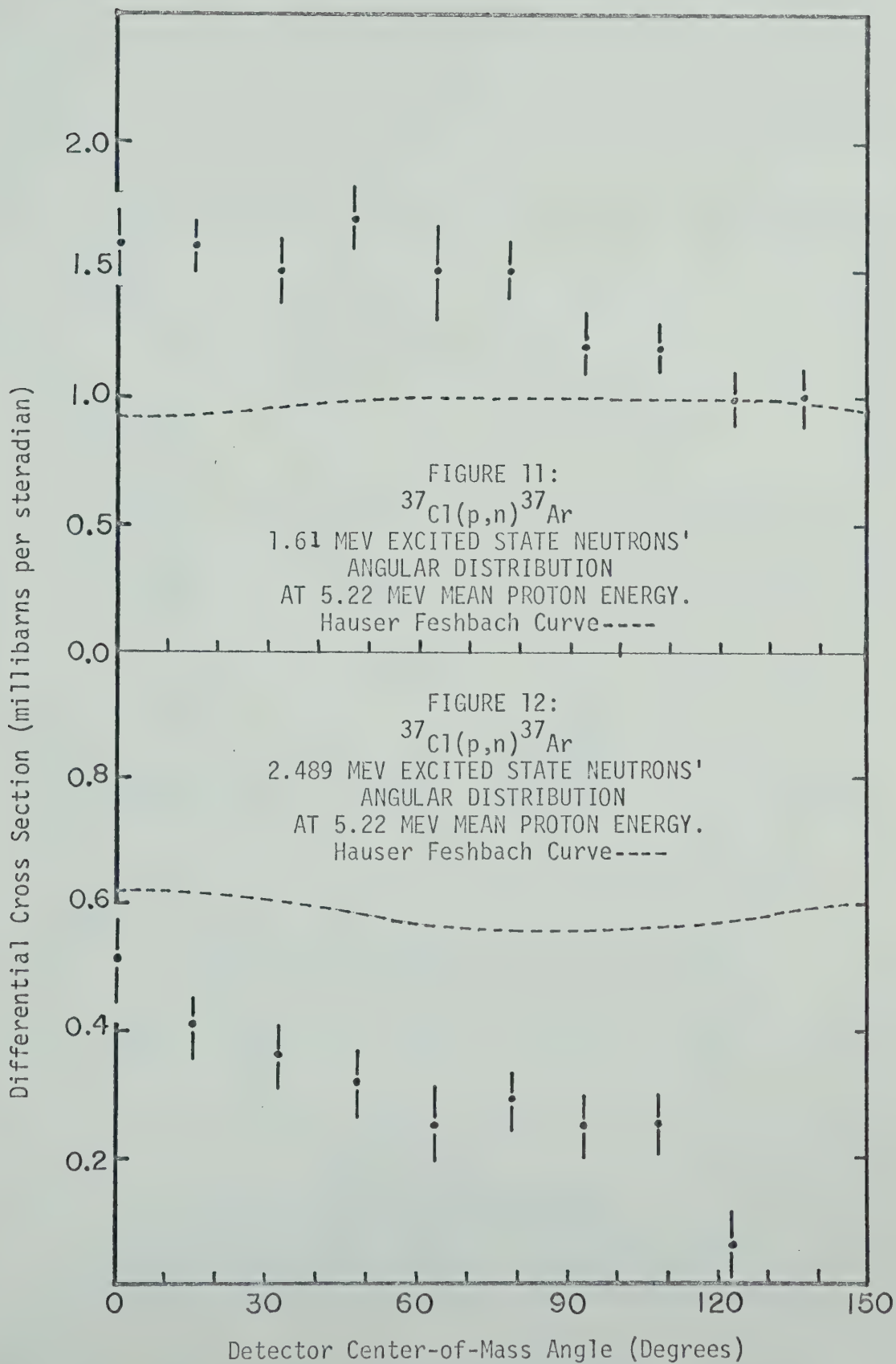
$$(26) \quad \text{S.D}(\bar{x}) = \frac{1}{n} \sqrt{\sum_i \epsilon_i^2}$$

where x_i and ϵ_i are the cross section and random error for each detector at each angle. The resulting neutron angular distributions are given in figures 9-13.

4.3 Random Errors

There were two significant random errors contributing to the ϵ_i . The first was due to the statistics of proton interaction with the nuclei of the target and assumed a Poisson distribution. Consequently it was taken to be $\sqrt{N}$ as an absolute error or $1/\sqrt{N}$ as a relative error,





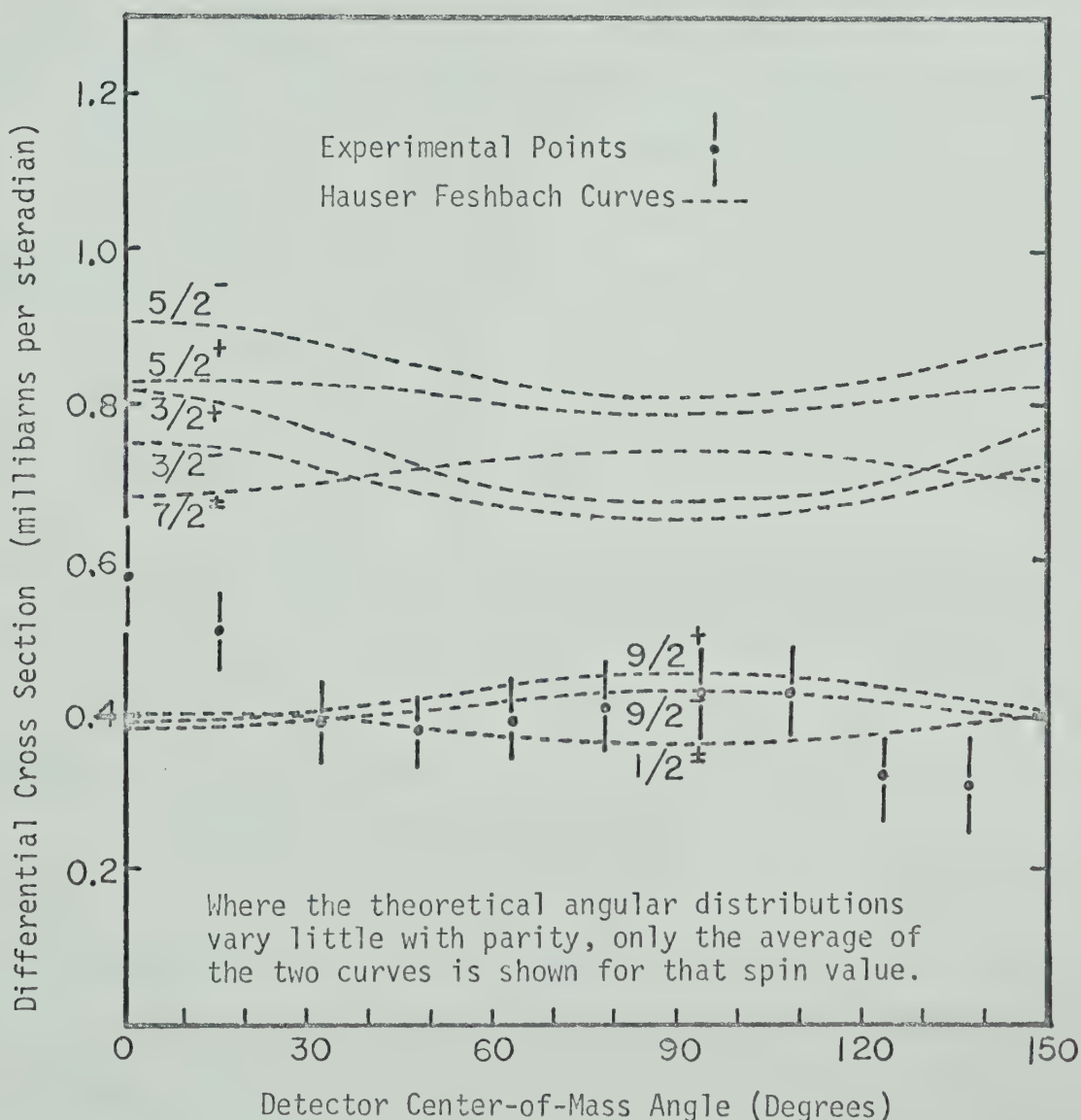


FIGURE 13:
 $^{37}\text{Cl}(p,n)^{37}\text{Ar}$
 2.22 MEV EXCITED STATE NEUTRONS'
 ANGULAR DISTRIBUTION AT 5.22 MEV MEAN PROTON ENERGY
 AND HAUSER FESHBACH CURVES, ASSUMING DIFF
 -ERENT SPINS AND PARITIES FOR THE 2.22 MEV LEVEL OF ^{37}Ar .

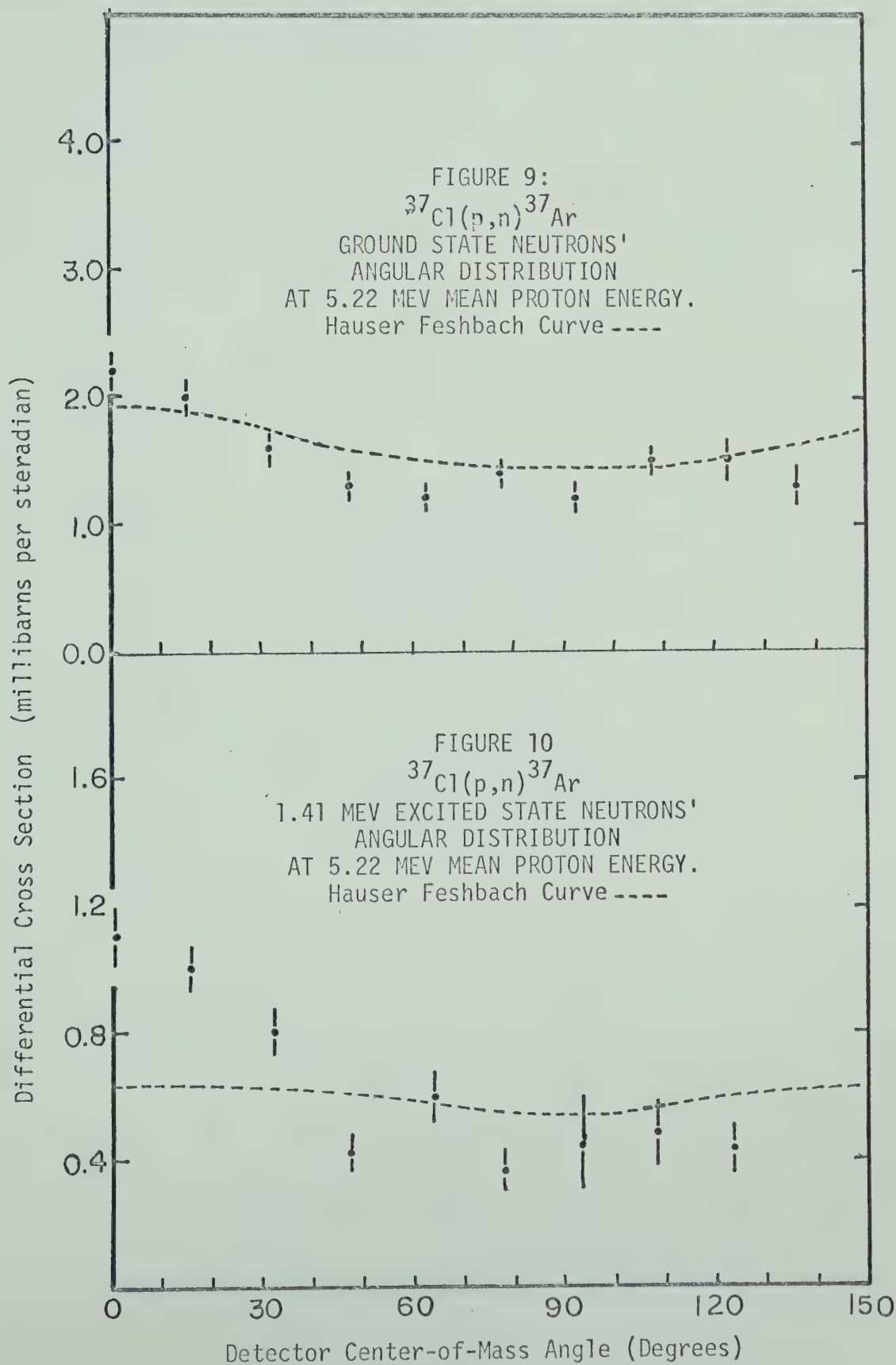
where N is the number of events.

The charge error discussed in the last section can be written as the relative error $S.D.(Q)/Q$. $S.D.(Q)$ denotes the standard deviation of the distribution of the charge, Q . This distribution was taken to have constant mean and a normal shape, though in reality, it was probably asymmetric. It then follows from equation 24 that the total relative error is given by taking the errors in quadrature:

$$(27) \quad (\epsilon_i / \frac{d\sigma}{d\Omega})^2 = \frac{1}{N} + (S.D.(Q)/Q)^2$$

where $1/\sqrt{N}$ has been substituted for the relative error in the number of counts, N . Thus the calculation of error in a given differential cross section measurement, ϵ_i , depended on an estimate for the relative error in Q .

The method of estimating this error in charge was a kind of iterative technique. That is, an average relative error in N was assumed and compared with the scatter of results from one detector about the mean of the results from both detectors. Then from equation 27 an estimate of $S.D.(Q)/Q$ could be calculated. The angular distributions of the neutrons from the ground state of ^{37}Ar was used, since this level had the lowest relative error in N and a constant efficiency error at all angles. The relative error in N was assumed to be the square root of the mean



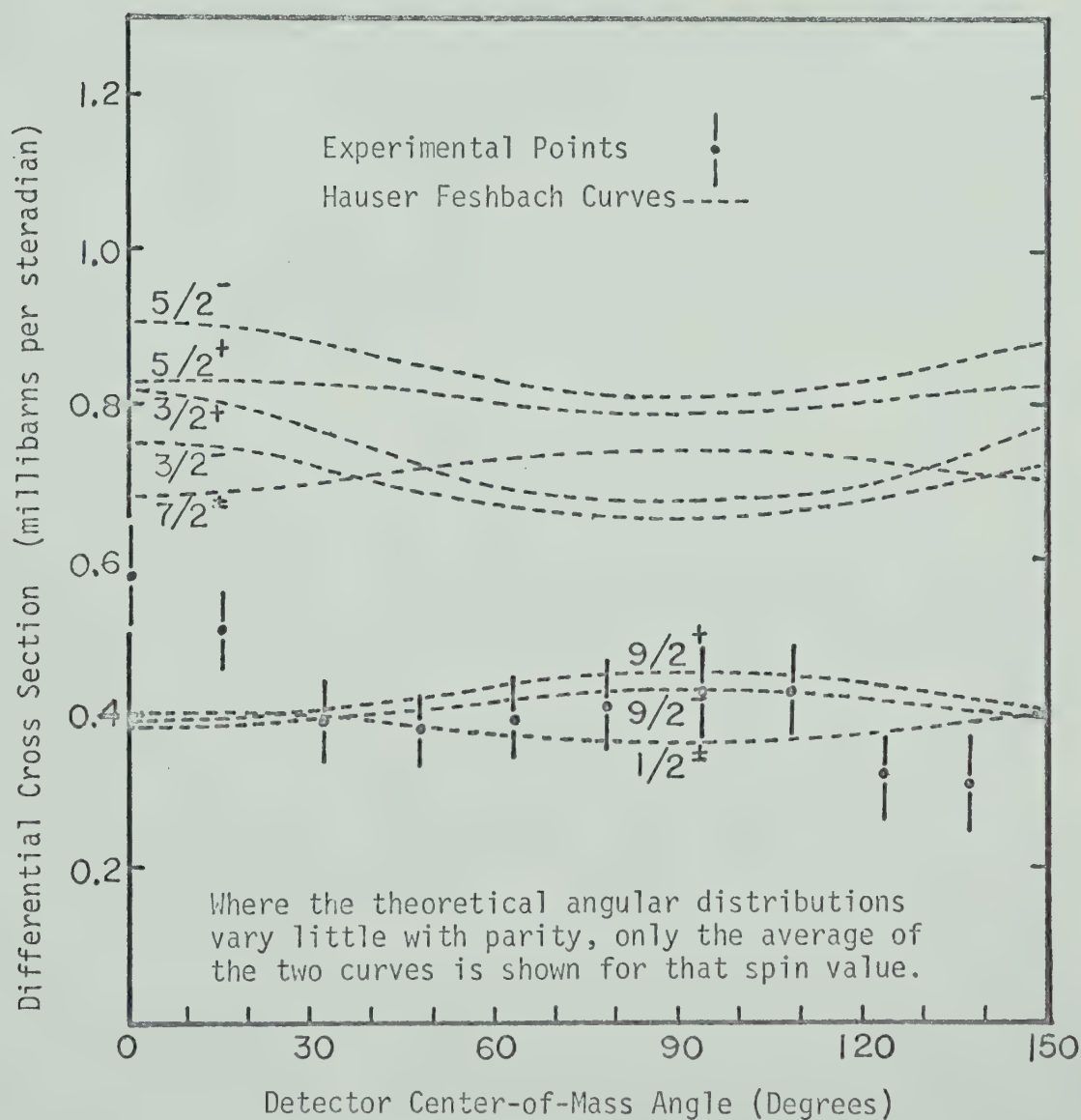


FIGURE 13:
 $^{37}\text{Cl}(p,n)^{37}\text{Ar}$
 2.22 MEV EXCITED STATE NEUTRONS'
 ANGULAR DISTRIBUTION AT 5.22 MEV MEAN PROTON ENERGY
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of the variances of N over all angles for this calculation. Similarly, the relative error in cross section $S.D.(\frac{d\sigma}{d\Omega})/\frac{d\sigma}{d\Omega}$ was taken to be the standard deviation of the quantity:

$$\left(\frac{\frac{d\sigma}{d\Omega} \text{ from detector \#1}}{\text{weighted average } \frac{d\sigma}{d\Omega} \text{ over both detectors}} - 1 \right)$$

over all the angles for which data was taken. In this case, the weighting was inversely proportional to the variance in the cross section due to the error in N only. To prohibit any systematic error in the efficiency of the two detectors artificially inflating this standard deviation, the angular distribution from detector #3 was normalized by least squares to the angular distribution of detector #1 before the weighted average in this expression was calculated. (This kind of normalization only pertains to this error calculation. It is not found in the mean cross section calculated in section 4.2). The resulting estimate of $S.D.(Q)/Q$ is 10% which is an average of two such values computed from two different sets of angles. If there are any other random errors that contribute to the scatter of the results of one detector about the other, they will be included in this figure.

Using the above value for $S.D.(Q)/Q$ and using, this time, the error in N calculated for a given differential cross section, the total error ϵ_i for each angle i was

found from the same equation 27. The error in the cross section for the given angle averaged over all the detectors, was then calculated from equation 26.

4.4 Systematic Errors

There were four important systematic errors. The first was concerned only with the neutrons from the 1.41 MeV and 1.61 MeV levels of Argon-37. Because the detectors were placed close to the target (3.75 meters) to increase the counting rate, the accompanying decrease in energy resolution resulted in partial overlapping of these two close levels. The double peak was analyzed into two peaks by least square methods using the program Tepel⁽⁶²⁾. This code fits the double peak with two peaks of arbitrary height and width but predetermined shape -- the reference shape, taken to be a smoothed version of the ground state peak.

A systematic error occurred in this calculation for results from detector #1 only. This was because single levels in its spectrum had small satellite peaks caused by poor electronics. The reference peak used was given this small satellite, but the result failed to adequately sort out the two levels. Consequently, a correction was made which minimized the mean square of the deviation of the detector #1 points from the detector #3 points. The

correction was a transfer of cross section from one level to the other, taken as a constant for all angles. The figures 10 and 11 are the result of this corrected Tepel sorting, averaged over the detectors.

The second error involved the calculation of detector efficiency. This follows from the uncertainty in determining the detector threshold as in section 3.4. It was a systematic error, slowly diverging from the true cross section until the neutron energy approaches the threshold when it could be quite large. An efficiency error also arose from the uncertainty in the calculation of detector efficiency, given a certain threshold. The angular distributions for all the levels except the 2.489 MeV level, could be systematically varied $\pm 10\%$ due to this efficiency error. This error for the 2.489 MeV level was $\pm 11\text{-}14\%$ for $\theta = 0^\circ$ to 125° respectively, because the neutrons from this level had energies close to the thresholds of the detectors. The latter estimate of error could quite conceivably be twice as much.

Finally there was the small anisotropy in neutron attenuation with varying neutron angle. This arose from neutron absorption in the cell walls and was only a few percent at most angles. Because of the difficulty in its calculation in the general case plus the large error in the charge described above, the figures were not corrected

for this effect. This error attained a maximum at the largest angle used, 135° , because at this angle the neutrons from the front part of the cell had to pass through the lip at the front of the cell mount on their way to the detector. (See figure 4). The ratio of attenuation at this angle to the minimum attenuation (which was at 90°) was 1.07.

4.5 Theoretical Cross Sections

The theoretical angular distributions were computed using a program "Hauser" written by W.R. Smith⁽⁶³⁾ and modified by N. Davison⁽⁶⁴⁾. Different values of spin and parity were assumed for the 2.22 MeV level of Argon-37 and the resulting angular distribution of neutrons from this state compared with experiment. This channel was only one of a large number available to the compound nucleus for decay. The known values of spin and parity were used for the other available levels of Argon-37. However, this data was not complete for the levels available in the $^{34}\text{S}+\alpha$ and $^{37}\text{Cl}+p$ channels. Consequently it was assumed that all levels above the ground state of ^{37}Cl and above the third excited state of ^{34}S were part of a continuum for which $N(E_x)$, the number of levels below the excitation energy E_x , could be obtained by fitting the straight line:

$$\log_{10} N(E_x) = a + b \sqrt{E_x}$$

to a log plot of the states given in the literature⁽⁶⁵⁾. Then a spin distribution of the form of equation 21 was assumed, with $\sigma^2 = 2.0 \sqrt{E_x}$. This value is consistent with that given by Von Witsch et al⁽⁴⁷⁾ for ^{38}Ar in section 2.2. The energy dependence of σ^2 suppressed the higher spins at lower excitation energies.

The transmission functions were calculated from the optical model parameters shown in table 1. The proton and neutron channels were based on parameters by Perey⁽⁶⁶⁾ and by Perey and Buck⁽⁶⁶⁾, respectively and gave the best theoretical fits to data of elastic scattering on various nuclei. If the neutron parameters of Rosen et al⁽⁶⁶⁾ had been used instead, only a 2% increase would have resulted in the cross section to ^{37}Ar states. If Rosen's proton parameters⁽⁶⁶⁾ had been used as well, the increase in theoretical cross section would have been from 4% to 13% depending on the ^{37}Ar level considered. In any event use of different optical model parameters did not materially alter any conclusions of this work. The alpha channel was based on the Set II parameters of L'Ecuyer and St. Pierre⁽⁶⁷⁾ which gave the best theoretical fit to elastic scattering of α on various nuclei as well as to the $(^3\text{He}, \alpha)$ reaction on several titanium isotopes. If their Set I parameters had been used, as well as those of Rosen et al⁽⁶⁶⁾ for the neutron and proton channels, the results would be 2% to 8% larger than those obtained using the optical model parameters of table 1.

TABLE 1
OPTICAL MODEL PARAMETERS

		Protons	Neutrons	Alphas
V_0	MeV	57.53	48	130
W_1	MeV	0	0	24
W_2	MeV	13.5	9.6	0
V_{so}	MeV	7.5	7.2	0
c_0	MeV/MeV	-0.55	-0.29	0
c_1	MeV/MeV	0	0	0
R	F.	1.25	1.25	1.25
R_0	F.	1.25	1.27	1.47
R_1	F.	1.25	1.27	1.47
a_0	F.	0.65	0.66	0.51
a_1	F.	0.47	0.47	0.51

4.6 Results and Conclusions

The lack of symmetry about 90° in figures 9-13 suggested interference took place between the compound nucleus mechanism and a smaller direct interaction mechanism. A broad maximum in the direct interaction angular distribution of neutrons from the first excited state of ^{37}Ar could occur at 0° or 67° for $\ell = 0$ or 1 respectively, where ℓ is the orbital angular momentum transferred to the residual nucleus. Of these, only $\ell = 0$ is permitted since the parity is unchanged in the reaction. Such a forward peak is seen in figure 10. A similar maximum in the angular distribution from the second excited state could occur at 0° , 67° , or 123° for $\ell = 0$, 1 or 2, respectively, of which only $\ell = 1$ is allowed by the selection rules with this reaction. Again a weak hump at about 67° can be found in figure 11. For the third excited state, a maximum could occur at 0° , 70° , or 134° for $\ell = 0$, 1 or 2, respectively. The forward peak of figure 13 suggests that a direct interaction interference component with $\ell = 0$ existed in the reaction. The work of A.J. Howard et al⁽²⁹⁾ has limited the spin of the third excited state to $(5/2, 7/2)$. For an assumed third excited state spin and parity of $5/2^+$, $\ell = 0, 2$ or 4 are allowed by the selection rules. For $7/2^+$, $\ell = 2, 4$, or 6 are permitted. With both $5/2^-$ and $7/2^-$, $\ell = 1, 3$, or 5 are

allowed for this reaction. Consequently figure 13 is consistent with the presence in this reaction of some direct interaction component and a $5/2^+$ spin and parity for the 2.22 MeV state of ^{37}Ar . Alternatively, some of these maxima could be examples of the resonance-type fluctuations.*

On the other hand the theoretical compound nucleus cross sections of neutrons from states of ^{37}Ar were several times larger than the experimental results. This could have been partly due to underestimating the number of excited states of ^{34}S and ^{37}Cl to which the compound nucleus could decay, states which would compete with decay to levels of ^{37}Ar . Another explanation is that ^{37}Ar and/or ^{37}Cl possess very small imaginary potential depths compared with the value found by Perey and others⁽⁶⁶⁾ for neighboring nuclides. Such an effect has been postulated by Bernard et al⁽⁵¹⁾ (see section 2.3) for this reaction.

Consequently the theoretical cross sections have been normalized by least squares to the ground state of ^{37}Ar , the level with the least experimental error and whose spin is well known. The normalization factor is 0.20. The resulting normalized theoretical angular distributions can be seen along with the experimental data in figures 9-13. The 1.6 MeV level of ^{37}Ar , which was separated from the 1.4 MeV level by the 'Tepel' code (section 4.4) does not show a good fit in spite of the fact that its spin is also well known. A poor fit is true

* See section 2.3.

of the 2.489 MeV ^{37}Ar level as well. In the case of the 2.22 MeV level of ^{37}Ar , to which it had been hoped a unique value of spin could be assigned by this method, figure 13 gives plausible fits for $1/2^{\pm}$ or $9/2^{\pm}$. However, this is at variance with the limits of $5/2$ and $7/2$ previously determined for this state. It appears that normalization to the ground state does not assure correct normalization of higher levels. Nor are the total reaction cross sections proportional to the number of magnetic substates of the residual nucleus, the ' $2I+1$ ' rule (with the exception discussed in section 4.1). Thus it would appear that the statistical theory is not completely adequate to describe the $^{37}\text{Cl}(p,n)^{37}\text{Ar}$ reaction at a mean proton energy of 5.22 MeV.

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APPENDIX I

TRANSMISSION FUNCTIONS AND FLUCTUATIONS

We will now show that the sum of transmission functions over all exit channels can be expressed in terms of compound nucleus level widths and densities. The form of the giant resonances affecting the average cross section through equation 20 can be derived from the Breit Wigner single-level amplitudes.

$$(28) \quad U_{nm} \approx e^{i(\Omega_m + \Omega_{m'})} (\Gamma_m^{1/2} \Gamma_n^{1/2}) / (E_0 - E - \frac{i}{2} \Gamma) \quad m \neq n.$$

This expression follows⁽²⁾ from expanding the logarithmic derivative of the external wave function at R' in a series about the resonance energy E_0 and retaining only the first two terms. The Γ is the resonance full width at half maximum and the Γ_n the partial width of channel n . Ω_n is the hard sphere potential scattering phase shift of channel n .

The dense overlapping series of resonances of widths Γ_λ^J and spacing D^J which underly the cross section fluctuations can also be given the Breit-Wigner form. So we write:

$$(29) \quad U_{nm} \approx e^{i(\Omega_m + \Omega_n)} \sum_{\lambda J} [(\Gamma_{\lambda m}^J)^{1/2} (\Gamma_{\lambda n}^J)^{1/2}] / (E_\lambda^J - E - \frac{i}{2} \Gamma_\lambda^J)$$

where $n \neq m$ and the pair (n, m) must be of the same J .

This expression can be strictly derived⁽¹¹⁾ from reaction theory when the partial widths (in contrast to the total widths) are much smaller than the level spacings. The Ericson fluctuations in the cross section are due to interference of different resonance terms in $|U_{nm}|^2$.

Let us now average $|U_{nm}|^2$ over an integral $\delta \gg \Gamma_\lambda^J$ or D^J . Here also, we call on the random phase approximation to ignore average cross terms between different resonances. The result is:

$$(30) \quad \overline{|U_{nm}|^2} = \overline{\sum_{\lambda}^J |(\Gamma_{\lambda m}^J)^{\frac{1}{2}}(\Gamma_{\lambda n}^J)^{\frac{1}{2}} / (E_\lambda^J - E - \frac{i}{2} \Gamma_\lambda^J)|^2}$$

$$= \overline{\sum_{\lambda}^J (\Gamma_{\lambda m}^J \Gamma_{\lambda n}^J) / ((E_\lambda^J - E)^2 + \frac{1}{4}(\Gamma_\lambda^J)^2)}.$$

Because

$$\frac{1}{\delta} \int_{-\infty}^{\infty} dE / ((E_\lambda^J - E)^2 + \frac{1}{4}(\Gamma_\lambda^J)^2) = 2\pi / \delta \Gamma_\lambda^J$$

and since there will be δ/D^J such integrals, we have

$$(31) \quad \overline{|U_{nm}|^2} = \frac{2\pi}{D^J \langle \Gamma_\lambda^J \rangle} \langle \Gamma_{\lambda m}^J \rangle \langle \Gamma_{\lambda n}^J \rangle \quad n \neq m$$

where we now use widths typical of that energy interval.

Comparing equation 31 with equation 15, equation 32 identifies:

$$(32) \quad T_m = 2\pi \langle \Gamma_{\lambda m}^J \rangle / D^J \text{ and } \sum_{cs''\ell''} T_p = 2\pi \langle \Gamma_\lambda^J \rangle / D^J.$$

This gives us the ratio of the level width to level spacings in terms of the transmission coefficient. That is, the compound nucleus parameters must be chosen to fit the gross variation of the transmission function over any giant resonances.

An analysis of the fluctuations themselves leads to a measure of $\langle \Gamma_\lambda \rangle$ so that, in principle, substitution in equation 32 leads to the level density alone if we assume the behavior of Γ_λ^J with J .

APPENDIX II

THE GAS CELL

A2.1 Construction and Use

Figure 4 shows the two portions of the brass cell: the cell itself backed by an air cooling compartment, and the mounted window. The window was 'epoxied'* to a bevelled brass mount, so shaped to allow free flexing of the window when the cell was pressured up. This mount was then sealed to the rest of the cell along the flanges by an 'O' ring. Exposed parts were covered with gold shielding to eliminate background radiation from the brass. The gas inlet line penetrated the base of the evacuated target chamber and the flow was controlled by valves on the gas handling apparatus. The pressure was measured with a mercury barometer. This cell had to be sealed during the experiment due to heating of the gas. A thermocouple had been installed to register cell temperature but was not working at the time of the experiment.

The difficulty in the construction of the cell had been in getting a window, thin enough to minimize beam energy loss and straggling, yet thick enough to form a seal between the gas target and the evacuated

* Epo-Tek #410 Part A and Epo-Tek 410-E Part B (hardener), Epoxy Technology Inc., 65 Grove St., Watertown, Mass. 02172.

target chamber. Both 0.11 thou inches of molybdenum foil and 0.125 thou inches of tantalum foil had been successfully used. The pressure differential across the window had to be in the direction of the mount or the window would unseat. In addition, the window could only withstand a fraction of an atmosphere pressure across it (i.e. the cells were tested to hold 30 cm Hg against a vacuum) and any change in pressure had to be very slow. This meant that some thought had to be given to pumping resistance in the relevant lines during the crucial stages of filling and emptying the cell. The target chamber and cell had to be pumped up (or down) alternately in stages.

The heating by the beam could have deteriorated the 'epoxy' seal. The heat flux from the window (about 0.3 watts with a 5 MeV proton beam and 2 μ amps current) encountered the highest thermal resistance at that point. This gradient had been reduced by using a silver based 'epoxy' cement. In this way cells have been run for at least 24 hours under a 10 watt beam. In view of the above comments, pretested spare mounted windows are mandatory for an experiment. For a discussion on the effect of the gas cell on energy resolution and isotropy of beam attenuation with angle, see sections 3.7 and 4.4 respectively.

A2.2 Recipe for Window Construction

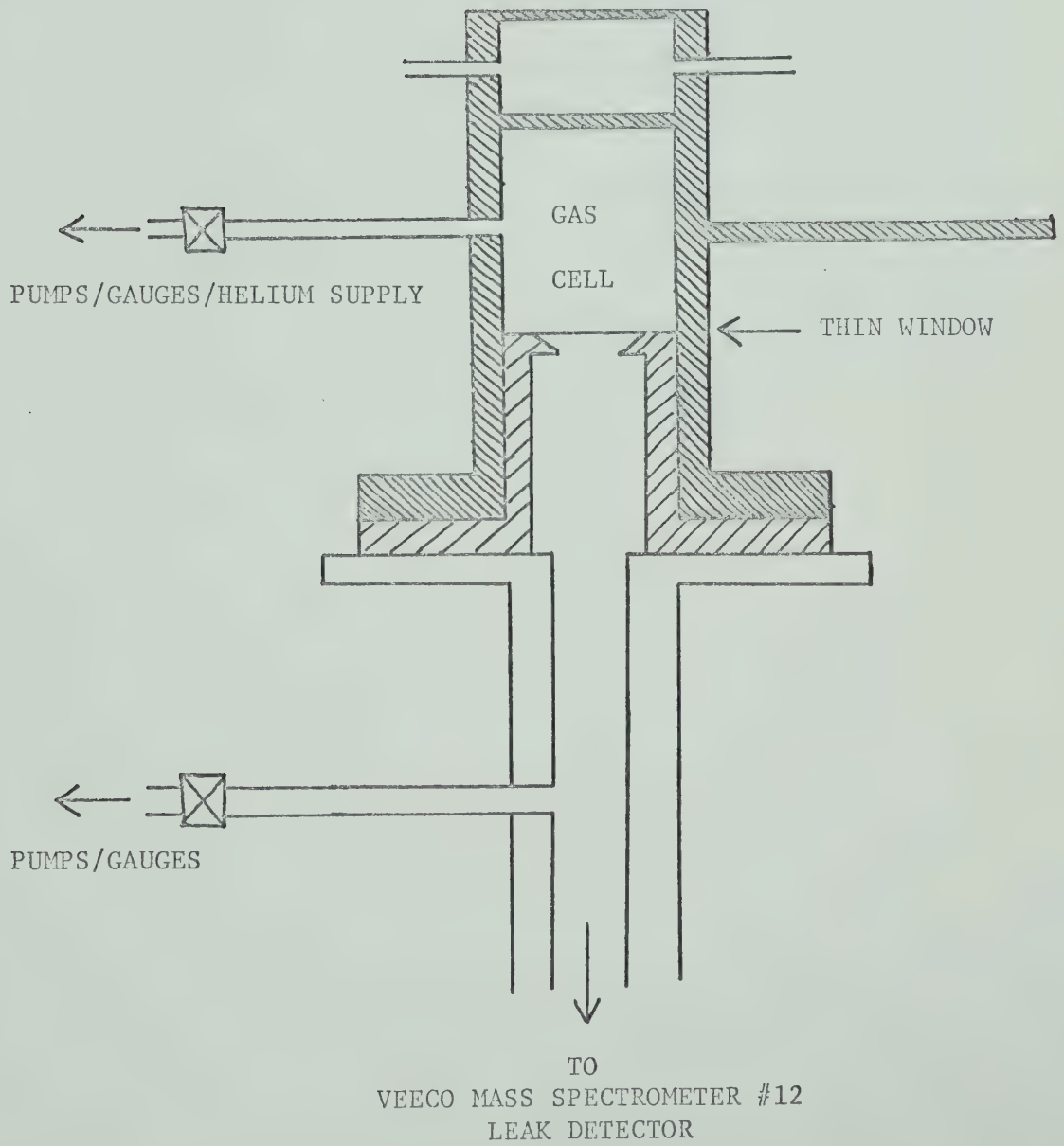
Select a square of pinhole free and fold-free foil, examined against a bright light. Then mix the hardener to the base, 1:15 by weight. Apply the 'epoxy' paste sparingly and evenly to the rim of a clean window mount. The mount is then inverted and gently pressed against the foil square. A smooth flat tool is used to work out small wrinkles and ensure a smooth seal all around the rim. This procedure should be done very slowly and with great care.

The assembly is then placed in an oven. A clean flat weight is set on the window to maintain the seal during heating. It was found necessary to separate the window and weight with a piece of Teflon tape to inhibit any sticking of the two that might take place. This could arise from cement squeezing around the edge of the foil, for instance. The assembly is heated at 120°C for $3\frac{1}{2}$ hours. The fragile mounted window is carefully removed and allowed to cool.

Since the window cannot withstand a one atmosphere pressure differential across it, it must be tested in the manner of figure 14. The Veeco* mass spectrometer leak detector senses extremely small helium leaks. The valves on the left are used for roughing down the system. The

* Veeco Vacuum Corp.

FIGURE 14:
GAS TARGET
LEAK TEST



comments regarding pressure changes given above apply here as well. The flow through the lower valve should always be away from the lower compartment to prevent any traces of helium reaching this part of the assembly. The lower section is then highly evacuated with the Leak Detector diffusion pump (open valves very slowly). The upper compartment is filled with helium to the desired pressure, keeping well below the rupture pressure of the window. In this way, very small leaks are immediately detected. With careful consideration of pressure changes, the assembly is brought up to air in stages.

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